2007 ELECTROWEAK INTERACTIONS AND UNIFIED THEORIES
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XLIIInd Rencontres de Moriond
La Thuile, Val d'Aoste, Italy - March 10 - 17, 2007

2007 Electroweak Interactions and Unified Theories

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Series: Electroweak Interactions and Unified Theories
La Thuile, Val d’Aoste, Italy March 10 - 17, 2007

2007 ELECTROWEAK INTERACTIONS AND UNIFIED THEORIES

edited by

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Jean Trần Thanh Vân
The Electroweak Session of the XLIIInd Rencontres de Moriond

was organized by

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with the active collaboration of:

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L. Okun (Moscow),
A. Pich (Valencia),
S. Pokorski (Warsaw),
G. Unal (Orsay),
D. Wood (Boston).
The XLIIInd Rencontres de Moriond were held in La Thuile, Vallée d'Aoste, Italie.
The first meeting took place at Moriond in the French Alps in 1966. There, experimental
as well as theoretical physicists not only shared their scientific preoccupations, but also
the household chores. The participants in the first meeting were mainly French
physicists interested in electromagnetic interactions. In subsequent years, a session on
high energy strong interactions was also added.

The main purpose of these meetings is to discuss recent developments in contemporary
physics and also to promote effective collaboration between experimentalists and
theorists in the field of elementary particle physics. By bringing together a relatively
small number of participants, the meeting helps to develop better human relations as
well as a more thorough and detailed discussion of the contributions.

This concern of research and experimentation of new channels of communication and
dialogue which from the start animated the Moriond meetings, inspired us to organize
a simultaneous meeting of biologists on Cell Differentiation (1980) and to create the
Moriond Astrophysics Meeting (1981). In the same spirit, we have started a new series
on Condensed Matter Physics in January 1994. Common meetings between biologists,
astrophysicists, condensed matter physicists and high energy physicists are organized
to study the implications of the advances of one field into the others. I hope that these
conferences and lively discussions may give birth to new analytical methods or new
mathematical languages.

At the XLIIInd Rencontres de Moriond in 2007, three Physics sessions, and one
Astrophysics session were held:

* March 10 - 17 "Electroweak Interactions and Unified Theories"
* March 11 - 18 "Gravitational Waves and Experimental Gravity"
* March 17 - 24 "QCD and High Energy Hadronic Interactions"
* March 18 - 24 "Venus Express Science Workshop"
I thank the organizers of the XLIIInd Rencontres de Moriond:


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Kevin Baines, Ludwik Celnikier, Pierre Drossart, David Grinspoon, Håkan Svedhem, Frederik Taylor, Dmitri Titov, Oliviere Witasse; for the "Venus Express Science Workshop" session

and the conference secretariat: P. Chemali, I. Cossin, G. Dreneau, G. Hérent, M. Joguet, Pham Duy Tu, F. Warin and V. Zorica.

I am also grateful to Enrico Belli, Ida Liseno, Monica Pelillo, Eric Agostini and Mirco Decci of the Planibel hotel who contributed through their hospitality and cooperation to the well-being of the participants, enabling them to work in a relaxed atmosphere.

These Rencontres were sponsored by the European Union "Marie Curie Conferences and Training Courses" Activity, the Centre National de la Recherche Scientifique (INSU, SPM and FP), the "Centre National d'Études Spatiales", the European Space Agency, the Institut National de Physique Nucléaire et de Physique des Particules (IN2P3-CNRS), the Commissariat à l'Energie Atomique (DAPNIA) and the National Science Foundation. I would like to express my thanks for their encouraging support.

I sincerely wish that a fruitful exchange and an efficient collaboration between the physicists and the astrophysicists will arise from these Rencontres as from the previous ones.

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I - SM Model parameters and Brout-Englert-Higgs Mechanism
Rencontres de Moriond 2007
FIRST RUN II MEASUREMENT OF THE W BOSON MASS WITH CDF

OLIVER STELZER-CHILTON for the CDF Collaboration
University of Oxford, Dept. of Physics, Denys Wilkinson Building,
Keble Road, OX1 3RH, Oxford, United Kingdom

The CDF collaboration has analyzed ~200 pb$^{-1}$ of Tevatron Run II data taken between February 2002 and September 2003 to measure the $W$ boson mass. With a sample of 63964 $W \rightarrow e\nu$ decays and 51128 $W \rightarrow \mu\nu$ decays, we measure $M_W = 80413\pm34(^{stat})\pm34(^{syst})$ MeV/c$^2$. The total measurement uncertainty of 48 MeV/c$^2$ makes this result the most precise single measurement of the $W$ boson mass to date.

1 Introduction

The $W$ boson mass is an important Standard Model (SM) parameter. It receives self-energy corrections due to vacuum fluctuations involving virtual particles. Thus the $W$ boson mass probes the particle spectrum in nature, including particles that have yet to be observed directly. The hypothetical particle of most immediate interest is the Higgs boson. The interaction of the other SM particles, in particular the $W$ and $Z$ gauge bosons, with the Higgs field is thought to impart mass to the SM particles. Thus the Higgs boson plays a critical role in the SM and it is very interesting to obtain a Higgs mass prediction. The $W$ boson mass can be calculated at tree level using the three precise measurements of the $Z$ boson mass, the Fermi coupling $G_F$ and the electromagnetic coupling $\alpha_{em}$. In order to extract information on new particles, we need to account for the radiative corrections to $M_W$ due to the dominant top-bottom quark loop diagrams. For fixed values of other inputs, the current uncertainty on the top quark mass measurement 170.9+1.8 GeV/c$^2$ corresponds to an uncertainty in its $W$ boson mass correction of 11 MeV/c$^2$. Measurements of the $W$ boson mass from Run I of the Tevatron and LEP with uncertainties of 59 MeV/c$^2$ and 33 MeV/c$^2$ respectively, yield a world average of 80392±29 MeV/c$^2$. It is clearly profitable to reduce the $W$ boson mass uncertainty further as a means of constraining the Higgs boson mass.
2 Measurement Strategy

At the Tevatron, W bosons are mainly produced by valance quark-antiquark annihilation, with initial state gluon radiation (ISR) generating a typical transverse boost of $O(10 \text{ GeV})$. The transverse momentum ($p_T^f$) distribution of the decay lepton has a characteristic Jacobian edge whose location, while sensitive to the W boson mass, is smeared by the transverse boost of the W boson. The neutrino transverse momentum ($p_T^\nu$) can be inferred by imposing $p_T$ balance in the event. The transverse mass, defined as $m_T = \sqrt{2p_T^\nu p_T^f (1 - \cos (\phi^\nu - \phi^f))}$, includes both measurable quantities in the W decay and provides the most precise quantity to measure the W boson mass. We use the $m_T$, $p_T^\nu$ and $p_T^f$ distributions to extract $M_W$. These distributions do not lend themselves to analytic parameterizations, which leads us to use a Monte Carlo simulation to predict their shape as a function of $M_W$. These lineshape predictions depend on a number of physical and detector effects, which we constrain from control samples or calculation. Important physical effects include internal QED radiation, the intrinsic W boson transverse momentum, and the proton parton distribution functions. Detector effects include external bremsstrahlung and ionization energy loss in the detector material, tracker momentum scale, calorimeter energy scale, resolution of the tracker and calorimeter, and the detector fiducial acceptance. In order to model and study these effects at the level of a part in $10^4$, we have developed a sophisticated, first-principles fast Monte Carlo simulation of the CDF detector. The W boson mass is extracted by performing a binned maximum-likelihood fit to the $m_T$, $p_T^\nu$ and $p_T^f$ distributions. We generate 800 templates as a function of $M_W$ between 80 GeV/$c^2$ and 81 GeV/$c^2$.

3 Momentum and Energy Scale Calibration

The key aspect of the measurement is the calibration of the lepton momentum. The charged lepton momentum is measured in a cylindrical drift chamber called the Central Outer Tracker (COT). The electron energy is measured using the central electromagnetic (EM) calorimeter and its angle measurement is provided by the COT trajectory. The COT track measurement sets the momentum scale for this analysis. The internal alignment of the COT is performed using high-$p_T$ cosmic rays that traverse diametrically the entire drift chamber. The momentum scale is set by measuring the $J/\Psi$ and $\Upsilon(1S)$ masses using the dimuon mass peaks. The $J/\Psi$ sample spans a range of muon $p_T$ (2-10 GeV/$c$), which allows us to tune our ionization energy loss model such that the measured mass is independent of muon $p_T$. We obtain consistent calibrations from the $J/\Psi$, $\Upsilon(1S)$ mass fits shown in Fig. 1 (left). The momentum scale extracted from the $Z \rightarrow \mu\mu$

![Diagram](image.png)

Figure 1: Left: Momentum scale summary: $\Delta p/p$ vs $1/p_T$ for $J/\Psi$, $\Upsilon(1S)$ and $Z$ boson dimuon samples. The dotted line represents the independent uncertainty between $J/\Psi$ and $\Upsilon(1S)$. Error bars are statistical only.
Right: Energy scale calibration using $E/p$ distribution from $W \rightarrow e\nu$ events.
mass fit, shown in the same figure, is also consistent, albeit with a larger, statistics-dominated uncertainty. The tracker resolution is tuned on the observed width of the $\Upsilon(1S)$ and $Z$ boson mass peaks.

Given the tracker momentum calibration, we fit the peak of the $E/p$ distribution of the signal electrons in the $W \rightarrow e\nu$ sample (Fig. 1 right) in order to calibrate the energy measurement of the electromagnetic (EM) calorimeter. The energy scale is adjusted such that the fit to the peak returns unity. The model for radiative energy loss is constrained by comparing the number of events in the radiative tail of the $E/p$ distribution. The calorimeter energy calibration is performed in bins of electron $E_T$ to constrain the calorimeter non-linearity. The calibration yields a $Z \rightarrow ee$ mass measurement of $M_Z = 91190.6 \pm 67_{\text{stat}} \text{ MeV}/c^2$, in very good agreement with the world average $(91187.6 \pm 2.1 \text{ MeV}/c^2)$; we obtain the most precise calorimeter calibration by combining the results from the $E/p$ method and the $Z \rightarrow ee$ mass measurement. The EM calorimeter resolution model is tuned on the widths of the $E/p$ peak and the $Z \rightarrow ee$ mass peak, separately for non-radiative and radiative electrons.

4 Hadronic Recoil Calibration

All particles recoiling against the $W$ or $Z$ boson are collectively referred to as the recoil. The recoil is computed as the vector sum of transverse energy over all calorimeter towers, where the towers associated with the leptons are explicitly removed from the calculation. The response of the calorimeter to the recoil is described by a response function which scales the true recoil magnitude to simulate the measured magnitude. The hadronic resolution receives contributions from ISR jets and the underlying event. The latter is independent of the boson transverse momentum and modeled using minimum bias data. The recoil response and resolution parameterizations are tuned on the mean and rms of the $p_T$-imbalance in $Z \rightarrow ll$ events as a function of boson $p_T$. We define the $\eta$ axis to be the geometric bisector of the two leptons and the $\xi$ axis to be perpendicular to $\eta$. We project the vector $p_T$-balance onto the $\eta$ and $\xi$ axes and compare the data distribution to the simulation. Fig. 2 (left) shows the mean of the $p_T$-balancing in $Z \rightarrow ee$ events as a function of $Z$ boson $p_T$. The quantity $p_T^Z$ is computed from the EM clusters of the

![Figure 2](image_url)

**Figure 2:** Left: Mean of the $p_T$-balancing as a function of $p_T$ in $Z \rightarrow ee$ events. Right: Projection of recoil along the lepton direction in $W \rightarrow \mu\nu$ events.

two $Z$ boson decay electrons and $u_{\eta}$ is computed from the recoil vector in the calorimeter. We cross-check the recoil model using $W$ and $Z$ boson data which show good agreement and validate the model. A very sensitive quantity to cross-check the recoil model in $W$ boson events is the projection of the recoil along the lepton direction ($u_{||}$). Fig. 2 (right) shows $u_{||}$ for $W \rightarrow \mu\nu$ events. We find good agreement between the data and the simulation.
5 Event Generation

We generate $W$ and $Z$ events with RESBOS$^4$, which captures the QCD physics and models the $W$ $p_T$ spectrum. RESBOS calculates the quintuple differential production cross section $\frac{d^5 \sigma}{dy \, dq_T \, dz_1}$, where $Q, y$ and $q_T$ are the boson invariant mass, rapidity, and transverse momentum respectively, and $dz_1$ is the solid angle element in the decay lepton direction. The RESBOS parametrization of the non-perturbative form factor is tuned on the dilepton $p_T$ distribution in the $Z$ boson sample. Photons radiated off the final-state leptons (FSR) are generated according to WGRAD$^5$. The FSR photon energies are increased by 10% (with an absolute uncertainty of 5%) to account for additional energy loss due to two-photon radiation$^6$. WGRAD is also used to estimate the uncertainty due to QED radiation from initial state (ISR) and interference between ISR and FSR. We use the CTEQ6M$^7$ set of parton distribution functions at NLO and apply their uncertainties to evaluate the systematic uncertainty on the $W$ boson mass.

6 Backgrounds

Backgrounds passing the event selection have different kinematic distributions from the $W$ signal and are included in the template fit according to their normalizations. Backgrounds arise in the $W$ boson samples from misidentified jets containing high-$p_T$ tracks and EM clusters, $Z \rightarrow ll$ where one of the leptons is not reconstructed and mimics a neutrino, $W \rightarrow \tau \nu$, kaon and pion decays in flight (DIF), and cosmic ray muons. The latter two are backgrounds in the muon channel only. Jet, DIF, and cosmic ray backgrounds are estimated from the data to be together less than 0.5%. The $W \rightarrow \tau \nu$ background is 0.9% for both channels, and the $Z \rightarrow ll$ is 6.6% (0.24%) in the muon (electron) channel, as estimated from Monte Carlo samples generated with PYTHIA$^8$ and a detailed GEANT-based detector simulation.

7 Results and Conclusions

The fits to the three kinematic distributions $m_T$, $p_T^f$ and $p_T^e$ in the electron and muon channels give the $W$ boson mass results shown in Table 1.

<table>
<thead>
<tr>
<th>Distribution</th>
<th>$W$ boson mass (MeV/c$^2$)</th>
<th>$\chi^2$/dof</th>
</tr>
</thead>
<tbody>
<tr>
<td>$m_T(e, \nu)$</td>
<td>80493±48(stat)±39(syst)</td>
<td>86/48</td>
</tr>
<tr>
<td>$p_T^f(e)$</td>
<td>80451±58(stat)±45(syst)</td>
<td>63/62</td>
</tr>
<tr>
<td>$p_T^e(e)$</td>
<td>80473±57(stat)±54(syst)</td>
<td>63/62</td>
</tr>
<tr>
<td>$m_T(\mu, \nu)$</td>
<td>80349±54(stat)±27(syst)</td>
<td>59/48</td>
</tr>
<tr>
<td>$p_T^f(\mu)$</td>
<td>80321±66(stat)±40(syst)</td>
<td>72/62</td>
</tr>
<tr>
<td>$p_T^e(\mu)$</td>
<td>80396±66(stat)±46(syst)</td>
<td>44/62</td>
</tr>
</tbody>
</table>

The transverse mass fit for the $W \rightarrow ev$ channel is shown in Fig. 3 (left) and for the $W \rightarrow \mu \nu$ channel in Fig. 3 (right). The uncertainties for the $m_T$ fits in both channels are summarized in Table 2. We combine the six $W$ boson mass fits including all correlations to obtain $M_W=80413±34$(stat)$±34$(syst) MeV/c$^2$. With a total uncertainty of 48 MeV/c$^2$, this measurement is the most precise single measurement to date. Inclusion of this result increases
the world average $W$ boson mass to $M_W = 80398 \pm 25$ MeV/c$^2$, reducing its uncertainty by 15%. The updated world average impacts the global precision electroweak fits, reducing the preferred Higgs boson mass fit by 6 GeV/c$^2$ to $M_H = 76^{+33}_{-23}$ GeV/c$^2$. The resulting 95% CL upper limit on the Higgs mass is 144 GeV/c$^2$ (182 GeV/c$^2$) with the LEP II direct limit included (excluded)\textsuperscript{3,9}. The direction of this change has interesting theoretical implications: as Fig 4\textsuperscript{10} shows, the $M_W$ vs $M_{top}$ ellipse moves a little deeper into the light-Higgs region excluded by LEP II and into the region favored by the minimal supersymmetry model (MSSM). While this is a one-sigma effect, it arouses further interest in higher precision measurements of $M_W$ (and $M_{top}$). Most of the systematic uncertainties in this measurement (Table 2) are limited by the statistics of the calibration samples used. Further improvements in the detector model and the production and decay model (e.g. QED radiative corrections) are likely to shrink the other systematic
Table 2: Systematic and total uncertainties for the transverse mass fits. The last column shows the correlated uncertainties, the last row is the quadrature sum of statistical and systematic uncertainty.

<table>
<thead>
<tr>
<th>Systematic (MeV/c²)</th>
<th>$W \rightarrow e\nu$</th>
<th>$W \rightarrow \mu\nu$</th>
<th>Common</th>
</tr>
</thead>
<tbody>
<tr>
<td>Lepton Energy Scale</td>
<td>30</td>
<td>17</td>
<td>17</td>
</tr>
<tr>
<td>Lepton Energy Resolution</td>
<td>9</td>
<td>3</td>
<td>0</td>
</tr>
<tr>
<td>Recoil Energy Scale</td>
<td>9</td>
<td>9</td>
<td>9</td>
</tr>
<tr>
<td>Recoil Energy Resolution</td>
<td>7</td>
<td>7</td>
<td>7</td>
</tr>
<tr>
<td>Selection Bias</td>
<td>3</td>
<td>1</td>
<td>0</td>
</tr>
<tr>
<td>Lepton Removal</td>
<td>8</td>
<td>5</td>
<td>5</td>
</tr>
<tr>
<td>Backgrounds</td>
<td>8</td>
<td>9</td>
<td>0</td>
</tr>
<tr>
<td>$p_T(W)$ Model</td>
<td>3</td>
<td>3</td>
<td>3</td>
</tr>
<tr>
<td>Parton Distributions</td>
<td>11</td>
<td>11</td>
<td>11</td>
</tr>
<tr>
<td>QED radiation</td>
<td>11</td>
<td>12</td>
<td>11</td>
</tr>
<tr>
<td>Total Systematics</td>
<td>39</td>
<td>27</td>
<td>26</td>
</tr>
<tr>
<td>Total Uncertainty</td>
<td>62</td>
<td>60</td>
<td>26</td>
</tr>
</tbody>
</table>

... uncertainties as well. CDF has now accumulated an integrated luminosity of about 2 fb⁻¹ and we look forward to a $W$ boson mass measurement with precision better than the current world average of 25 MeV/c², with the dataset already in hand.

Acknowledgments

I would like to thank my colleagues from the CDF collaboration in particular the $W$ boson mass group for their hard work on this important analysis. Sincere thanks also to the conference organizers and participants for a superb conference. This work was supported by the European Commission under the Marie Curie Programme.

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Diboson physics from the Tevatron

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Experimental studies of diboson production at the Tevatron provide stringent tests of the
electroweak gauge structure in the Standard Model. In these proceedings, we review the
latest results from the CDF and DØ experiments based on up to 1 fb$^{-1}$ of Run II data. The
discussion covers the well-established WW signal, the recent observation of WZ production
and the latest searches for ZZ. Measurements of the W$\gamma$ production cross section and a study
of the charge-signed rapidity difference, as well as measurements of the Z$\gamma$ production cross
section are also discussed.

1 Introduction

The non-abelian SU(2)$_L \otimes$ U(1)$_Y$ structure of the Standard Model (SM) implies the existence of
non-zero self-couplings of the gauge bosons. The SM makes detailed predictions for the
structure of the corresponding trilinear gauge couplings (TGC). Experimental studies of diboson
production provide a stringent test of the SM predictions: measurements of total and differential
cross sections are sensitive to anomalous couplings (AC). ACs would in general lead to increased
cross sections, in particular for large transverse momenta of the gauge bosons.

In these proceedings, we will give a brief overview of diboson results from Run II of the CDF
and DØ experiments at the Fermilab Tevatron. The results from the Tevatron, a $p\bar{p}$ collider
at $\sqrt{s} = 1.96$ TeV, are complementary to those from LEP, since the Tevatron probes different
combinations of TGCs at higher $\hat{s}$, where New Physics effects might manifest themselves. As of
early 2007, CDF and DØ have recorded more than 2 fb$^{-1}$ of data each, and their datasets are
expected to increase to 4-8 fb$^{-1}$ per experiment over the next two years. The analyses presented
here are based on up to the first fb$^{-1}$. 
2 WW production

WW production receives contributions from the WWZ and WW$\gamma$ couplings. The cross section is predicted \(7\), at next-to-leading order, to be \(12.4 \pm 0.8 \) pb. This channel is also an important background in searches for a possible high-mass Higgs boson in \(H \to WW^*\), i.e. has to be understood before \(H \to WW^*\) can be observed.

The WW signal in Run II is well established. Despite its small branching fraction, the case of two semileptonic W decays is particularly well suited in a hadron collider environment. CDF and DØ require two isolated high-$p_T$ leptons (ee, e$\mu$, $\mu\mu$) of opposite charge plus missing transverse energy. Same-flavour combinations consistent with the $Z$ mass are rejected. The measured cross sections are

\[
\sigma(p\bar{p} \to W^+W^-) = 13.6 \pm 2.3 \text{ (stat)} \pm 1.6 \text{ (syst)} \pm 1.2 \text{ (lumi)} \text{ pb} \quad \text{CDF prel.}^{1,7}, \quad 825 \text{ pb}^{-1}, \\
\sigma(p\bar{p} \to W^+W^-) = 14.6^{+5.8}_{-5.1} \text{ (stat)}^{+1.8}_{-3.0} \text{ (syst)} \pm 0.9 \text{ (lumi)} \text{ pb} \quad \text{DØ}^{7}, \quad 224-252 \text{ pb}^{-1},
\]

in good agreement with the SM prediction.

3 WZ production

In contrast to WW, WZ production is sensitive to only the WWZ vertex. WZ production is unavailable at LEP, i.e. studies of WZ at the Tevatron allow the first study of WWZ coupling without assumptions on WW$\gamma$ coupling. The cross section is predicted \(8\) to be \(3.7 \pm 0.1 \) pb. DØ and CDF use the clean $WZ \to \ell\nu\ell'^-\ell^-$ modes, where $\ell$ is an electron or muon. Including contributions from $W$ and/or $Z$ decays to $\tau$ leptons, the branching fraction of WZ to three $e$ or $\mu$ leptons is 1.8%. A pair of isolated high-$p_T$ leptons consistent with the $Z$ mass is required, plus a third high-$p_T$ lepton and missing transverse energy. In the DØ analysis, all three lepton candidates must pass lepton identification requirements based on the calorimeter and the muon detectors. The CDF analysis also considers, in addition, candidate events where up to two candidate leptons are reconstructed only as isolated central tracks that are not fiducial to calorimeters or muon detectors. The dominant backgrounds are $Z/\gamma^* +$ jets (fake isolated lepton candidate) and ZZ production. The measured cross sections $\sigma(p\bar{p} \to WZ)$ are summarised below:

<table>
<thead>
<tr>
<th>Luminosity</th>
<th>Measured</th>
<th>Signal significance</th>
<th>Predicted $^2$</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF prel. $^7$</td>
<td>1.1 fb$^{-1}$</td>
<td>$5.0^{+1.3}_{-1.1} \text{ (stat)} \pm 0.4 \text{ (syst)} \text{ pb}$</td>
<td>$6.0\sigma$</td>
</tr>
<tr>
<td>DØ prel. $^7$</td>
<td>760 - 880 pb$^{-1}$</td>
<td>$4.0_{-1.5}^{+1.4} \text{ (stat + syst) pb}$</td>
<td>$3.3\sigma$</td>
</tr>
</tbody>
</table>

For example the CDF result is based on 16 candidate events with an expected background of $2.7 \pm 0.28 \text{ (stat)} \pm 0.33 \text{ (syst)} \pm 0.09 \text{ (lumi)}$. The dilepton invariant mass for the same-flavour opposite-charge dilepton pair closest to the $Z$ mass in the CDF candidate events is shown in Fig. ???. WZ production has now been observed and the measured cross sections are consistent with the SM.

4 ZZ production

The SM does not predict any self-coupling of $Z$ bosons, i.e. the vertices $ZZ\gamma$ and $ZZZ$ are forbidden. The leading SM diagram for $(Z/\gamma^*)\,(Z/\gamma^*)$ production at the Tevatron is shown in Fig. ???. The predicted $^2$ $p\bar{p} \to ZZ$ cross section at the Tevatron, including one-loop corrections, is 1.60 pb. At this conference, a new search $^7$ for ZZ by DØ based on 1 fb$^{-1}$ of data has been reported for the first time. This search uses the clean leptonic $eeee$, $ee\mu\mu$ and $\mu\mu\mu\mu$.
modes. Candidate events are retained if both $Z/\gamma^*$ boson candidates have an invariant mass above 30 GeV in at least one of the possible lepton pairings. Due to the small number of expected $Z/\gamma^*$ pair events and the sensitivity to single lepton identification cuts, considerable effort was spent optimizing the acceptance and lepton efficiencies for the three channels. The expected event yields, as well as the number of observed events are summarised in Tab. 2. The signal acceptance is dominated by geometric acceptance and identification inefficiencies for the leptons. The small expected backgrounds are dominated by $t\bar{t}$ events, single $Z$ bosons produced with two jets, and $Z$ bosons produced with a photon and a jet. The observed yield is consistent with the background expectation, and DØ set a limit of 4.3 pb at 95% confidence level on $p\bar{p} \rightarrow (Z/\gamma^*)/(Z/\gamma^*)$ for an invariant mass requirement on both $Z/\gamma^*$ bosons of 30 GeV or greater. This result is consistent with the CDF limit of 3.8 pb at 95% confidence level based on leptonic final states. Both limits are also consistent with the SM.

Just after this conference, CDF have reported a new complementary search for $ZZ$ in the $t^+t^-\nu\nu$ final state. In that search, contributions from the different modes (e.g. $WW$, $ZZ$, Higgs) that contribute to this final state are disentangled using event probabilities calculated from all measured kinematic information and leading order differential cross sections. Run II is expected to provide enough data for an observation of $ZZ$, assuming the SM cross section.

5 $W\gamma$ production

$W\gamma$ production can occur by the radiation of a photon in the initial or final state, or via the $WW\gamma$ vertex. Anomalous $WW\gamma$ couplings would lead to an enhancement in the production cross section and enhance the rates for events with a high-$E_T$ photon or high transverse mass of the $W\gamma$ system. Photon identification at hadron colliders is challenging due to significant
backgrounds from jets with, e.g., a leading $p_T$. For this analysis, CDF use central ($|\eta| < 1.1$) photon candidates above $E_T > 7$ GeV; DØ use the same $E_T$ threshold and a wider angular coverage ($|\eta| < 1.1$ and $1.5 < |\eta| < 2.5$). Both analyses use clean semi-leptonic $W$ decays and require the photon to be well-separated from the lepton ($\Delta R(l, \gamma) > 0.7$ in $\eta/\phi$ space). The two analyses yield the following cross sections, which are in good agreement with SM expectations:

\begin{align*}
\text{DØ prel.}^{7} (0.9 \text{ fb}^{-1}) & : \quad E_T(\gamma) > 7 \text{ GeV}, \Delta R(l, \gamma) > 0.7, M_{T(l, \gamma)} > 90 \text{ GeV} \\
\text{muon channel} \quad \sigma(p\bar{p} \rightarrow l\nu\gamma X) &= 3.21 \pm 0.49 \text{ (stat + syst)} \pm 0.20 \text{ (lumi)} \text{ pb} \\
\text{electron channel} \quad \sigma(p\bar{p} \rightarrow l\nu\gamma X) &= 3.12 \pm 0.40 \text{ (stat + syst)} \pm 0.19 \text{ (lumi)} \text{ pb} \\
\text{theory}^{2} \quad \sigma(p\bar{p} \rightarrow l\nu\gamma X) &= 3.21 \pm 0.08 \text{ (PDF)} \text{ pb} \\
\text{CDF prel.}^{2} (1 \text{ fb}^{-1}) & : \quad E_T(\gamma) > 7 \text{ GeV}, \Delta R(l, \gamma) > 0.7 \\
\text{muon channel} \quad \sigma(p\bar{p} \rightarrow l\nu\gamma X) &= 19.11 \pm 1.04 \text{ (stat)} \pm 2.40 \text{ (syst)} \pm 1.11 \text{ (lumi)} \text{ pb} \\
\text{theory} \quad \sigma(p\bar{p} \rightarrow l\nu\gamma X) &= 19.3 \pm 1.4 \text{ (syst)} \text{ pb}
\end{align*}

The kinematic distributions of the candidate events are also in agreement with SM predictions. As an example, the photon $E_T$ from the CDF analysis is shown in Fig. ???. DØ increase sensitivity to ACs by studying the charge-signed rapidity difference $Q \cdot \Delta \eta$, where $Q$ is the charge of the lepton and $\Delta \eta = \eta(\gamma) - \eta(l)$ is the rapidity difference between the photon and the lepton. The distribution of this quantity in background-subtracted DØ data is shown in Fig. ???, along with the expectation from an SM Monte Carlo simulation. In the absence of final state radiation and detector effects, the differential cross section is expected to be essentially zero at $Q \cdot \Delta \eta = 0$ (a manifestation of the “radiation amplitude zero” effect), but these effects are included in the MC simulation and the dip is somewhat washed out. The data are in agreement with the SM prediction ($\chi^2/\text{ndof} = 16/12$). For illustration, an alternative MC prediction based on one particular assumption on ACs (a choice that leads to zero magnetic dipole moment of the $W$ boson, and that is therefore expected to remove the radiation amplitude zero) is also shown in Fig. ???.

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*Footnotes:*

1. Charge-singlet rapidity difference
2. Monte Carlo predictions

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*References:*

2. DØ collaboration (2007), arXiv:0707.4906
6 $Z\gamma$ production

In the SM, photons do not directly couple to $Z$ bosons, i.e., the vertices $ZZ\gamma$ and $Z\gamma\gamma$ are forbidden. In this model, the $Z\gamma$ final state can, however, be produced by final state radiation (FSR) off both charged leptons from $Z/\gamma^* \rightarrow l^+l^-$ or by one of the partons in the $p$ or $\bar{p}$ through initial state radiation (ISR). Both CDF and DØ have measured the cross section for $Z\gamma$ production in datasets of 1 fb$^{-1}$ per experiment. In the CDF analysis, the $Z/\gamma^* \rightarrow e^+e^-$ channel is used, DØ consider both $Z/\gamma^* \rightarrow e^+e^-$ and $Z/\gamma^* \rightarrow \mu^+\mu^-$. The $Z/\gamma^*$ decay can be cleanly reconstructed by requiring two isolated high-$p_T$ leptons. The experimental challenge is again photon identification, and the dominant background is $Z/\gamma^* +$ jet production with the jet misidentified as photon. Both experiments require the photon to be central ($|\eta| < 1.1$), above $E_T > 7$ GeV and well separated from the leptons ($\Delta R(l,\gamma) > 0.7$ in $\eta/\phi$ space). Fig. ?? shows the distribution of the candidate events from the CDF analysis in the $M(e^+e^-)/M(e^+e^-)$ plane. The different production processes like FSR, ISR and Drell-Yan (DY) lead to distinctive patterns in the distribution. CDF and DØ measure cross sections for different cuts on $M(l^+l^-)$ and $M(l^+l^-$), and compare the measurements to SM theory predictions:

**CDF prel.** (1.1 fb$^{-1}$): $E_T(\gamma) > 7$ GeV, $\Delta R(e,\gamma) > 0.7$

$M(ee\gamma) > 40$ GeV: $\sigma(p\bar{p} \rightarrow Z\gamma) \cdot \text{BR}(Z \rightarrow ll) = 4.9 \pm 0.3 \text{(stat)} \pm 0.3 \text{(syst)} \pm 0.3 \text{(lumi)} \text{ pb}$

theory

$M(ee\gamma) > 100$ GeV: $\sigma(p\bar{p} \rightarrow Z\gamma) \cdot \text{BR}(Z \rightarrow ll) = 1.4 \pm 0.1 \text{(stat)} \pm 0.2 \text{(syst)} \pm 0.1 \text{(lumi)} \text{ pb}$

theory

**DØ** (1 fb$^{-1}$): $E_T(\gamma) > 7$ GeV, $\Delta R(l,\gamma) > 0.7$

$M(ll) > 30$ GeV: $\sigma(p\bar{p} \rightarrow Z\gamma) \cdot \text{BR}(Z \rightarrow ll) = 4.96 \pm 0.30 \text{(stat + syst)} \pm 0.30 \text{(lumi)} \text{ pb}$

theory

$\sigma(p\bar{p} \rightarrow Z\gamma) \cdot \text{BR}(Z \rightarrow ll) = 4.74 \pm 0.22 \text{(syst)} \text{ pb}$

The measurements are in agreement with theory expectations. The measured photon $E_T$ spectra, shown in Figs. ?? and ?? are also in agreement with predictions. DØ interpret their results in terms of limits on the parameters in a generalised parameterisation of anomalous $ZZ\gamma$ and $Z\gamma\gamma$ couplings in the framework of Ref. ?.
7 Conclusion

We have reviewed the latest diboson results from CDF and DØ based on up to 1 fb\(^{-1}\) of data. WW production has been established for the first time at a hadron collider early in Run II of the Tevatron. WZ production has recently been observed, and ZZ production should be established over the next year or two. W\(\gamma\) and Z\(\gamma\) production have also been studied. No deviations from the SM have been observed. The interpretation of some of these results in terms of limits on parameters that describe anomalous couplings still have to be finalised. At the same time, the next two years will tell us what 4–8 fb\(^{-1}\) per experiment have in store for us. These impressive datasets will allow even more precise tests of the SM.

Acknowledgments

It is a pleasure to thank my CDF and DØ colleagues for the exciting and fruitful collaboration, and our Tevatron colleagues for the excellent luminosity. I also wish to thank the organisers for the nice conference.

W AND Z PHYSICS AT TEVATRON

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Recent measurements of W and Z boson properties at the Tevatron proton anti-proton collider are presented. A b-specific jet energy scale using $Z \rightarrow b\bar{b}$ signal has been extracted from CDF Run II data. The leptonic decay of Z boson in pairs of $e^+e^-$ has been studied by both CDF and DØ experiments to measure the Z boson differential cross section. Leptonic decay of W boson has also been used, by DØ experiment, to measure the muon charge asymmetry.

1 Introduction

Studying the electroweak sector of the Standard Model (SM) is of major importance for the physic program of the CDF and DØ experiments for Run II at Tevatron. The large number of W and Z bosons, produced in this proton anti-proton collider operating at $\sqrt{s} = 1.96$ TeV, allow to measure and test precisely their properties. For instance differential cross section measurements of $Z/\gamma^*$ boson leptonic decay or the study of lepton charge asymmetry in W decays provide useful knowledge of the parton distribution functions. The hadronic decays of W and Z bosons are of great interest as well. For example the hadronic decay of W boson is used in top mass analysis as a way to constrain the jet energy scale. On the other hand a large-sized signal of $Z \rightarrow b\bar{b}$ decays allows a precise measurement of a b-quark specific jet energy scale, which is useful for all physic signal that involve b-jets.

We present here recent results on W and Z boson physics performed at CDF and DØ based on integrated luminosity ranging from 0.2 to 1 fb$^{-1}$.

2 Determination of a b-jet Energy Scale Using $Z \rightarrow b\bar{b}$ Decay

Extracting the $Z \rightarrow b\bar{b}$ resonance allows both a precise measurement of the energy scale of b-quark jets and a determination of the b-jet energy resolution. The reduction of the uncertainty in the b-jet energy scale helps all precision measurements of the top quark mass, while a determination of the b-jet energy resolution is important for the search of a low-mass Higgs boson. The signal, most notably, opens the doors to a direct test of algorithms that attempt to increase the resolution of the b-jet energy measurement. These algorithms are a critical ingredient for the observability of the Higgs boson at the Tevatron if $M_H < 135$ GeV/c$^2$.

The jet energy scale (JES) is a factor which measures the discrepancy between the effect of detector response and energy corrections in real and simulated hadronic jets. When dealing with b-jets, however, one has to cope with several peculiarities of their fragmentation and decay
properties. These effects have to be accurately modeled if one is to use a generic JES factor extracted from jets not containing heavy flavors to $b$-jets, such as those present in a $t\bar{t}$ decay.

This analysis is based on an integrated luminosity of $584 \text{ pb}^{-1}$ collected by the Collider Detector at Fermilab (CDF II) from March 2004 to February 2006. A specific trigger was designed as a mean of acquiring a sizable amount of $Z$ decays to pairs of $b$-jets, thereby providing a convenient calibration line for the energy measurement of $b$-jets.

A preliminary sample was selected by requesting events with two high transverse energy central jets ($E_T > 22 \text{ GeV}$ and detector pseudorapidity $|\eta_2| \leq 1.0$) and by requiring both leading jets to be taggable\(^a\). The number of data events surviving these preliminary requirements is $18,128,488$. When both leading jets are required to be identified as $b$-jets (i.e., “tagged”) by the presence of secondary decay vertices using SECVTX algorithm\(^7\), the number of selected events drops to $699,990$.

Although the previous set of cuts allows to select events with low dijet invariant mass spectrum, which is crucial to observe a signal peak distinct from the background mass turn-on, the collected sample is largely dominated by QCD dijet background. However, because of the differences between $Z \rightarrow b\bar{b}$ process and QCD dijet production, requesting back to back events with low extra-jet radiation helps enhance the signal fraction. In order to reduce the background and initial and final state radiation processes, additional cuts on the azimuthal difference between the leading two jets, $\Delta\phi_{12}$, and the transverse energy of the third jet, $E_T^3$ were requested ($\Delta\phi_{12} > 3.0$ and $E_T^3 < 15 \text{ GeV}$). The number of selected events surviving all cuts (including tagging of the leading two jets) is $267,246$ events.

The number of signal and background events, in selected data, and the $b$-jet energy scale factor, $k$, are measured simultaneously through an unbinned likelihood procedure. Pythia Monte-Carlo $Z \rightarrow b\bar{b}$ samples are used to construct a signal probability density function which has the dijet invariant mass and the $b$-jet energy scale as parameters. The mass shape of the background collected in the double tagged data, after the kinematical selection of events, is modeled using a data-driven method. A background enriched sample, orthogonal to the signal region, is selected and used to compute a tag rate shape, function of the dijet invariant mass. This ratio is then multiplied by the mass distribution of events with two taggable jets accepted by the kinematical cuts that define the signal region to construct the background template. Several background shapes can be constructed depending on the choice of the background enriched region. The background model that best fits the sideband is used in an unbinned likelihood fit over the total dijet mass spectrum to estimate the $b$-JES factor and number of events of signal. The others possible background models are used to calculate the systematic errors related to background modeling.

The probability that observed data is described as an admixture of background events and $Z \rightarrow b\bar{b}$ events with a data/MC $b$-jet energy scale $k$ is given by an unbinned likelihood function. The fit precision on $k$ is improved by including a gaussian constraint on the expected number of events of signal (estimated from Monte-Carlo). Figure ?? shows the result of the unbinned likelihood fit to double tagged data. The goodness of this fit is estimated calculating the $\chi^2/NDF$ and gives $104/75$. The final results, including all systematic calculations, are, for the $b$-jet energy scale factor:

$$k = 0.974 \pm 0.011^{+0.017}_{-0.014} (\text{syst.}) = 0.974^{+0.020}_{-0.028} (\text{stat + syst}),$$

and for the number of fitted signal events:

$$N_{\text{sig}} = 5674 \pm 448 (\text{stat.})^{+1473}_{-1250} (\text{syst.}) = 5674^{+1540}_{-1250} (\text{stat + syst}) \text{ events.}$$

The large $Z \rightarrow b\bar{b}$ decay signal extracted from CDF II data allowed a measurement of the $b$-jet energy scale with a total precision of $\%$. This result is relevant for $b$-jets in an energy range close to that of jets from $Z$ decay, and additional studies are needed to exploit it in different signals involving $b$-jets.

\(^a\)Le to contain at least two tracks measured by the silicon tracker.
Figure 1: Result of unbinned likelihood fit performed to double tagged dijets data (blue points). The data-driven background shape and Monte-Carlo signal p.d.f are shown respectively in green and red. The fit returns 5674 ± 448 events of signal and a b-JES of 0.974 ± 0.011. The inset on the upper right shows the data minus background distribution (blue points) and the signal shape (in red) normalized to the fitted number of events of signal.

3 $Z/\gamma^*$ Boson Rapidity Distribution

Kinematic distributions of $Z/\gamma^*$ bosons produced in hadronic collisions provide a wealth of informations on the fundamental interactions involved. Measurement of the rapidity distribution and total cross section of Drell-Yan pairs in the $Z/\gamma^*$ boson region provide a stringent test of QCD calculations in the leading order (LO) and next to leading order (NLO). At leading order, $Z/\gamma^*$ bosons are produced through the annihilation of a quark and an anti-quark, with the partons in the proton and anti-proton carrying momentum fractions $x_1$ and $x_2$, respectively. The momentum fractions $x_1(x_2)$ are related to the rapidity ($y$) of the $Z$ boson via the equation:

$$x_{1,2} = \frac{M_{Z/\gamma^*}e^{\pm y}}{\sqrt{s}}$$  \hspace{1cm} (1)

where $\sqrt{s}$ is the center of mass energy and $M_{Z/\gamma^*}$ the mass of the $Z/\gamma^*$ boson. Both CDF and DØ measured the differential cross section distribution, $\frac{d^2\sigma}{dy}$, using dielectrons from $Z/\gamma^*$ decays.

CDF used a dataset of 1.1 fb$^{-1}$ collected using inclusive single central electron and dielectron triggers. $Z$ bosons decays are selected as events with two high transverse energy electrons ($E_T > 20 - 25$ GeV). The dielectron data sample consists of three different topologies depending whether the electrons, from the $e^+e^-$ pair, fall in the central ($|\eta| < 1.1$) or the plug calorimeters ($1.2 < |\eta| < 2.8$). After the cuts a total of 91,362 electron pairs are selected. The electron trigger efficiencies as a function of their transverse energy $E_T$ are measured using the data. The overall trigger efficiency of the dielectron pairs versus rapidity is almost 100%. Geometric and kinematic acceptance are modeled using Pythia $Z \rightarrow e^+e^-$ event generator combined with GEANT simulation of the CDF detector. The total acceptance times efficiency is flat up to $y \sim 2.0$ and is non-zero up to $y = 2.9$.

The main source of background, for this analysis, is the QCD dijet background, which is estimated from data itself using the electron isolation energy$^b$ distribution. The background rates for electroweak processes ($WW, WZ, t\bar{t}$ inclusive, $W$ inclusive) are estimated using simulation.

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$^b$The isolation energy is defined as the energy contained within a cone in $(\eta, \phi)$ of $\Delta R \leq 0.4$ around the electron minus the energy of the electron itself.
The differential cross section of electron pairs is defined as:

$$\frac{d\sigma(Z/\gamma^*)(y)}{dy} = \frac{N_{obs}(y) - N_{bg}(y)}{A \times \epsilon(y) \int L dt}$$

where $N_{obs}(y) - N_{bg}(y)$ is the number of background-subtracted events, $A \times \epsilon(y)$ the combined acceptance and efficiency and $\int L dt$ the total integrated luminosity. Figure 2 shows the $d\sigma(Z/\gamma^*)/dy$ distribution as a function of $|y|$. In this measurement the largest systematic uncertainty is associated with the measurement of the silicon tracking efficiency. The differential cross section measurement was compared to theory predictions and the NNLO calculation with NLO CTEQ6.1 PDF was found to be most consistent with the data. The total measured $Z/\gamma^* \rightarrow e^+e^-$ cross section, from integrating $d\sigma(Z/\gamma^*)/dy$, is $\sigma = 265.9 \pm 1.0 \pm 1.1$ pb.

![Figure 2: Measured $d\sigma/dy$ for $p\bar{p} \rightarrow Z/\gamma^* \rightarrow e^+e^-$ in CDF II data. The crosses are the measurements and the solid line is the theory prediction.](image)

DO experiment also performed a measurement of $Z/\gamma^*$ boson rapidity distribution, using 0.4 fb$^{-1}$ of data. The normalized differential cross section was found to be consistent with MRST04 NNLO calculation.

4 Z Boson Transverse Momentum

A precision measurement of the Z boson transverse momentum provides a sensitive test of the theory for weak boson production, it also helps reduce the theoretical uncertainty on the precision W mass measurement. DO performed a preliminary measurement of inclusive differential cross section as a function of Z boson transverse momentum in an invariant mass range between 70 and 110 GeV/c$^2$ using 960 pb$^{-1}$ of collected data.

The selection criteria for candidate Z bosons decaying in $e^+e^-$ pair requires two isolated electromagnetic (EM) clusters in the calorimeters that pass electron identification criteria, have high transverse momenta, and an invariant mass consistent with that of the Z boson. A total of 63,901 events are selected and identified as $Z/\gamma^*$ candidates. Electron selection efficiencies are measured from the data, using the “tag and probe” method, and are parameterized in terms of the electron $E_T$ and detector pseudorapidity, $\eta_{det}$.

ResBos and PHOTOS are used as the $Z/\gamma^*$ event generator. A parameterized Monte-Carlo simulation is used to simulate the effect of detector smearing and to include the measured electron identification efficiencies. The electron energy scale and resolutions are tuned using the Z data.
As in the previously described analysis, the main source of background comes from QCD dijets. To estimate the size of this background, samples that contain mostly background events, are used to get the shape of the invariant mass distribution from this source. Other sources of background ($Z \rightarrow \tau \tau$ and dibosons) are studied using Monte-Carlo and are found to be negligible.

The measured Z boson $p_T$ spectrum is smeared due to detector resolution effects. To be able to compare with theory directly, the detector effects are first unfolded. The largest uncertainty on the unfolded Z $p_T$ spectrum arise from the dependence of the efficiencies of the lepton identification requirements on the boson transverse momentum. The $\frac{dN}{dp_T}$ distribution with both statistical and systematic uncertainties is shown in Figure 3. The data are compared to theoretical prediction and shown to be in good agreement.

![Unfolded Z boson $p_T$ distribution](image)

**Figure 3:** Unfolded Z boson $p_T$ distribution. The uncertainty contains both statistical and systematic uncertainties.

5 W Boson Charge Asymmetry

On average, the u-quark in the proton carries a larger fraction of the proton’s momentum than the d-quark. This implies that at production, a $W^{+}(-)$ boson is typically boosted in the $p(\bar{p})$ direction. This leads to a charge asymmetry, as a function of the W boson rapidity, $y_W$. It is difficult to measure the $W^{+}$ rapidity due to the fact that the longitudinal momentum of the neutrino from the W decay cannot be measured directly. However, the same information can be accessed by measuring the charge asymmetry of the W boson decay products. The lepton charge asymmetry is a convolution of the W production charge asymmetry and the asymmetry from the (V-A) decay. The lepton asymmetry can thus be used to probe the parton distributions. Both the DØ (muon channel) and CDF (electron channel) collaborations have measured the lepton charge asymmetry, defined as:

$$A(y) = \frac{N^+(y) - N^-(y)}{N^+(y) + N^-(y)}.$$  \hspace{1cm} (3)

The main experimental challenge is to keep the lepton-charge misidentification rate low (∼0.01% for DØ muon channel), and to understand possible charge-dependencies of selection efficiencies.

The DØ measurement is based on $\int L dt \approx 230 \text{ pb}^{-1}$ of recorded data. Good $W \rightarrow \mu \nu$ candidates are selected by requiring events with a high transverse momentum muon ($p_T > 20$ GeV), significant missing transverse energy ($E_T > 20$ GeV) and large W transverse mass ($M_T >$
40 GeV). After all selection cuts 189,697 W candidates are selected in the data. The largest source of contamination in the sample comes from electroweak backgrounds ($Z \to \mu\mu, W \to \tau\nu, Z \to \tau\tau$). These backgrounds are estimated using Monte-Carlo samples. The other major source of contamination comes from the multijet background (semi-leptonic decay) and is estimated from data.

The final result of the DØ measurement, is given in figure ???. Even though the uncertainties on the experimental points are dominated by the finite data-sample-size, some sensitivity to PDF’s is already obtained (assuming the W boson decays according to the standard model). Thus, with more data, this measurement will constrain the PDF’s.

![Figure 4: Measured and predicted muon charge asymmetry, as a function of muon rapidity. The theoretical curves are based on NLO calculations (RESBOS and CTEQ6.1M PDFs).](image)

6 Conclusions

A large signal of $Z \to b\bar{b}$ decays from Tevatron collider data has been successfully extracted. This signal allowed to perform the first precise measurement of a b-specific jet energy scale. This measurement is of interest for all signals involving b-jets (top quark mass measurements, Higgs boson searches).

On the other hand, precision measurements of differential W and Z boson production cross sections have been performed. The results agree well with Standard Model predictions. Measurements of lepton charge asymmetries as a function of lepton rapidity are already sensitive to PDF’s. In the near future, several times more data will be analyzed. This will provide a better determination of PDF’s, which in its turn is important for reducing systematic uncertainties on, for example, the W boson mass measurement.

References

Early Electroweak Measurements in CMS and ATLAS

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Electroweak processes will be extremely important during the first phases of the LHC programme. They will be used to calibrate the ATLAS and CMS detectors, to understand the underlying event environment and to obtain background control samples for a large variety of possible new-physics signals. A progressive scenario can be envisaged: large statistics samples of W and Z decays into leptons will be available for an integrated luminosity of 10 pb⁻¹. W and Z production accompanied by jets will be studied with 100 pb⁻¹, diboson signals like \( pp \to WZ \) will be observable with 150 pb⁻¹ and a detailed understanding of the detector performance using \( tt \) events could be possible with luminosities as low as 300 pb⁻¹. Exploiting such a rich scenario is one of the objectives and challenges for the first year of LHC operation.

1 Introduction

Electroweak processes are important at the LHC for many reasons. On one hand, they can be considered as "known physics". The electroweak sector of the Standard Model has been precisely studied at previous colliders, like LEP, HERA and Tevatron, and analyses at the Tevatron collider suggest that W, Z, Drell-Yan and top production at the LHC should be well under control. Electroweak reactions will be used as reference samples to calibrate the LHC detectors and understand their response to muons, electrons, photons and jets. They will also be used to quantify the performance of sophisticated tools like b-tagging.

On the other hand, electroweak processes are not so well known processes. LHC collisions will explore a new energy domain. Basic ingredients in any physics study are the parton density functions (PDFs), calculated for the corresponding parton fraction value and at the relevant energy scale of the process. Of particular importance at the LHC are the gluon PDFs, due to their dominance in many physics processes (Higgs and top production, for instance). At
present, this gluon component is one of the largest sources of PDF uncertainty at the LHC. Top physics is still in its early stages, and LHC should mark the start of a new phase of precision measurements for this sector.

We bring the attention to the fact that many new physics signals will have electroweak processes as dominant backgrounds. In many cases, the new physics signals will even contain Z, W, γ* (Drell-Yan) or top decays in the final state. Therefore it is extremely important to understand electroweak contributions before claiming any discovery. This could be the case even at the earliest stages of the LHC, since the average center-of-mass energy explored at a 14 TeV pp collider is one order of magnitude above that of previous accelerators (LEP, Tevatron). Figure 1 illustrates this point. Even at an integrated luminosity of 1 fb⁻¹ LHC will explore scales not accessible at Tevatron. The measurements in this region - with the present knowledge - will be dominated by uncertainties in the electroweak sector.

![Figure 1: Early new physics searches at CMS with Drell-Yan processes. Electroweak and PDF uncertainties are the dominant sources of systematic uncertainties at high invariant masses (CMS).](image)

2 Inclusive Z and W production

Decays of W and Z bosons into leptons will be the first electroweak signals to be exploited at the LHC. Thanks to their large cross sections (above 10 and 1 nb, respectively), thousands of W and Z leptonic decays will be already at our disposal for integrated luminosities as low as 1 pb⁻¹.

New studies on these channels can be found in the recently released Physics TDR³ of the CMS collaboration. The emphasis has been put in a start-up oriented selection of samples with high purity. The criteria are chosen to be minimally dependent on calibration uncertainties or possible imperfections in the detector response. Figure 2 illustrates the Z → μμ case. In order to avoid potential systematic biases due to an imperfect matching with the inner tracking system or inefficiencies in the muon system, one of the muons is identified as an isolated track with high transverse momentum (p_T). The high efficiency of the dimuon trigger in Z → μμ events can be used to estimate the single muon efficiency in W → μν events. Conversely, Z → μμ events fired only by the single muon trigger can be used to study possible regional inefficiencies in the muon trigger system. In summary, a loose approach at the selection level gives enough lever arm to study detector performances and inefficiencies in detail.

CMS assigns experimental systematic uncertainties of 1.1% and 2.2% for the pp → Z + X → μμ + X and pp → W + X → μν + X measured cross sections, respectively. The overall
Figure 2: Selection of $Z \rightarrow \mu\mu$ events in CMS. On the left, one can observe the better and uniform coverage obtained as a function of pseudo-rapidity by relaxing the selection criteria on one of the muons in the event. The plot on the right shows the level of trigger redundancy in the sample.

Uncertainty is larger (2.3% and 3.3%), due to theoretical uncertainties in the determination of the acceptance, and in particular in the description of the $p_t$ spectrum of the boson. In the absence of detailed studies and comparisons with data for new next-to-leading-order Monte Carlos (NLO) like MC@NLO, a systematic uncertainty of $\sim 2\%$ was assigned from the difference between the acceptances obtained with a leading-order Monte Carlo (HERWIG 4) and MC@NLO. The differences in the $p_t$ spectra of the selected muons are shown in Figure 3.

Figure 3: Muon $p_t$ spectrum for $W \rightarrow \mu\nu$ (left) and $Z \rightarrow \mu\mu$ (right) events. The differences in shape between the predictions of leading-order (LO) and next-to-leading-order (NLO) Monte Carlos is visible. All histograms are normalized to 1. In the CMS studies, the difference in acceptance between the LO and NLO cases is assigned as a theoretical uncertainty ($\sim 2\%$).

Another important source of theoretical uncertainties is the choice of PDFs. PDF uncertainties have a limited impact on the acceptance ($< 1\%$ for $|\eta| < 2.1$), and therefore on the experimental measurement. However, their effect on the absolute normalization of the signal is rather large, of order $5 - 7\%$. Let us note that these uncertainties can not be easily reduced, since they manifest as a global normalization factor, largely uncorrelated with variations in shape in the fiducial volume used. Unless more precise PDF sets (from HERA, for instance) become available in the next future, this will be a limiting factor in comparisons between experiment and theory, as well as in measurements of the luminosity via $W \rightarrow \mu\nu$ and $Z \rightarrow \mu\mu$ event counting.
3 W/Z plus jets measurements

Besides its intrinsic interest as a QCD measurement, W and Z production accompanied by jets will be a unique tool to understand the typical underlying jet-event structure at the LHC and to reduce jet energy scale uncertainties (via Z + jets). This type of events are also one of the dominant backgrounds in new physics searches with leptons in the final state, as well as in top precision studies. Compared with Tevatron, the cross sections for these channels are higher. On the negative side, semileptonic top decays constitute a significant background component for W + jets at the LHC.

The CMS collaboration has recently released a study of W/Z plus jets associated production for transverse jet energies $E_T > 50$ GeV. The basic ingredients of the analysis are: a) a consistent and robust definition of a jet in the selected events and b) very stringent lepton isolation criteria in order to suppress QCD backgrounds. The numbers of selected events for an integrated luminosity of 1 fb$^{-1}$ are shown in Tables 1 and 2. Let us note that, even with luminosities as low as 100 pb$^{-1}$, a significant measurement of the Z + 4 jets cross section is possible.

Table 1: Number of selected events and breakdown of the different background components in the W + jets analysis of CMS. An integrated luminosity of 1 fb$^{-1}$ is assumed. Top backgrounds are comparable of even larger than the signal for W + ≥ 3 jets.

<table>
<thead>
<tr>
<th>Channels</th>
<th>W ≥ 1 jet</th>
<th>W ≥ 2 jet</th>
<th>W ≥ 3 jet</th>
<th>W ≥ 4 jet</th>
</tr>
</thead>
<tbody>
<tr>
<td>W + jets (signal)</td>
<td>260652</td>
<td>56702</td>
<td>10964</td>
<td>2164</td>
</tr>
<tr>
<td>Z + jets</td>
<td>9340</td>
<td>3237</td>
<td>972</td>
<td>259</td>
</tr>
<tr>
<td>τt + jets</td>
<td>12897</td>
<td>11842</td>
<td>9052</td>
<td>5420</td>
</tr>
<tr>
<td>WW/WZ/ZZ + jets</td>
<td>1077</td>
<td>714</td>
<td>386</td>
<td>151</td>
</tr>
<tr>
<td>Total</td>
<td>283966</td>
<td>72495</td>
<td>21374</td>
<td>7994</td>
</tr>
</tbody>
</table>

Table 2: Number of selected events and breakdown of the different background components in the Z + jets analysis of CMS. An integrated luminosity of 1 fb$^{-1}$ is assumed. Compared to the W + jets case, the relative contributions of the different backgrounds, and in particular of τt, are negligible.

<table>
<thead>
<tr>
<th>Channels</th>
<th>W ≥ 1 jet</th>
<th>W ≥ 2 jet</th>
<th>W ≥ 3 jet</th>
<th>W ≥ 4 jet</th>
</tr>
</thead>
<tbody>
<tr>
<td>Z + jets (signal)</td>
<td>35109</td>
<td>6185</td>
<td>977</td>
<td>156</td>
</tr>
<tr>
<td>τt + jets</td>
<td>64</td>
<td>58</td>
<td>49</td>
<td>32</td>
</tr>
<tr>
<td>WW/WZ/ZZ + jets</td>
<td>33</td>
<td>17</td>
<td>5</td>
<td>2</td>
</tr>
<tr>
<td>Total</td>
<td>35206</td>
<td>6260</td>
<td>1031</td>
<td>190</td>
</tr>
</tbody>
</table>

4 Diboson production

Both ATLAS and CMS collaborations have carried out studies on diboson production with leptons in the final state. Even if one of main purposes of these studies is to evaluate the LHC potential to measure anomalous triple-gauge boson couplings, the main objective for the startup phase is to measure the diboson cross sections at these new energies. These processes are also the dominant backgrounds in critical new physics searches, like $H \rightarrow WW \rightarrow$ leptons or $H \rightarrow ZZ \rightarrow$ leptons.

Table 3 shows the preliminary ATLAS result for the WZ → leptons channel. Figure 4 shows the equivalent analysis reported by CMS in its TDR. Significant observations (at > 5σ level) of WZ and ZZ production are expected in CMS for integrated luminosities of 150 pb$^{-1}$ and 1 fb$^{-1}$, respectively.
Table 3: Number of selected signal and background events in a preliminary study of WZ production of the ATLAS collaboration. An integrated luminosity of 1 fb$^{-1}$ is assumed. Numbers for the different leptonic final states used in the analysis are also given.

<table>
<thead>
<tr>
<th>WZ production</th>
<th>$N_{WZ}$</th>
<th>$N_{Z_{ee}}$</th>
<th>$N_{Z_{rp}}$</th>
<th>$N_{Z_{pp}}$</th>
<th>$N_{Z_{tot}}$ (1 fb$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Signal</td>
<td>16.9</td>
<td>17.1</td>
<td>21.9</td>
<td>19.8</td>
<td>75.7</td>
</tr>
<tr>
<td>Background</td>
<td>1.71</td>
<td>0.88</td>
<td>1.73</td>
<td>2.00</td>
<td>6.32</td>
</tr>
<tr>
<td>Significance</td>
<td>12.9</td>
<td>18.2</td>
<td>16.7</td>
<td>14.0</td>
<td>30.1</td>
</tr>
</tbody>
</table>

Figure 4: $Z$ reconstructed mass obtained in CMS studies of WZ and ZZ production with leptons in the final state. Observations at the 5σ level are already possible with luminosities of 150 pb$^{-1}$ (WZ) and 1 fb$^{-1}$ (ZZ).

5 Top production

The predicted $pp \to t\bar{t}$ cross section at the LHC is huge (800 pb). It is two orders of magnitude larger than the cross section measured at the Tevatron, due to the dominance of the gluon production mechanism $pp \to gg \to t\bar{t}$. Understanding top production requires a good understanding of the response of the whole detector. Given the high statistics that will be available in the early phases, the reverse will become true at the LHC: relatively pure samples of $t\bar{t}$ events will be used to understand a wide spectrum of detector-related issues. Lepton identification and isolation, jet energy scales, missing transverse energy and b-tagging performance are some examples.

Both ATLAS and CMS collaboration are considering progressive scenarios as a function of the integrated luminosity $L$, which can be summarized as follows:

- $L = 20 - 30$ pb$^{-1}$: with this luminosity, studies similar to those currently under way at the Tevatron are possible. Focusing on final states containing leptons and measuring cross sections for the first time at the LHC would be the main objectives of this phase.

- $L = 200 - 300$ pb$^{-1}$: strategies and methods are established. More precise measurements of the cross sections and first measurements of the top mass should be available. This is the phase where detector effects will be studied in detail for the first time.

- $L \sim 1$ fb$^{-1}$: an optimal understanding of the detector response is expected. Full exploitation of the physics potential is the main objective. This phase should mark the beginning of a top-physics precision era, with accurate measurements of the top mass and couplings.
Figure 5 (ATLAS) shows an example of the accuracies (~1%) that can be reached in the calibration of the jet-energy scale using $t\bar{t}$ events\textsuperscript{9}. The energy scale is determined from a $\chi^2$ fit on a pure sample of $W$-bosons decaying into jets. The well known value of the $W$-mass is used as the main constraint.

Figure 5 also shows an example of b-tagging studies in CMS using $t\bar{t}$ events\textsuperscript{10}. B-enriched jet samples are obtained in $t\bar{t}$ events by applying stringent leptonic criteria (all lepton samples) or criteria minimally dependent on lifetime tracking information ("semileptonic" samples). This allows the determination of b-tagging efficiencies in data as a function of the jet energy and pseudorapidity with uncertainties better than 10% for luminosities of the order of 1 fb$^{-1}$.

Figure 5: Examples of detector-related studies using $t\bar{t}$ events with an integrated luminosity on 1 fb$^{-1}$. Left: jet energy calibration at the 1% level using $W$ mass constraints (ATLAS). Right: determination of b-tagging efficiencies with uncertainties better than 10% (CMS).

References

TOP MASS MEASUREMENT AT THE TEVATRON

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Results on the measurement of the top quark mass from the two Tevatron collider experiments, CDF and DØ are presented here. We start with an introduction to top quark physics and the techniques used by both experiments to achieve a precise measurement of the top quark mass. The most recent and/or most precise measurements performed by the two experiments in different final state channels are then described. The measurements are performed on datasets corresponding to integrated luminosities up to 1 fb$^{-1}$. We conclude by presenting the latest world average value for the mass of the top quark: $M_{\text{top}} = 176.9 \pm 1.1(\text{stat.}) \pm 1.5(\text{sys.})$.

1 Top Quark Physics and the Top Quark Mass

In proton anti-proton collisions at Tevatron energies, $\sqrt{s} = 1.96$ TeV, top quarks are primarily produced in pairs (a top and an anti-top quark) via strong interactions. We measure the top quark mass in this production mode. Approximately 85% of the production cross-section at Leading Order (LO) is initiated by $q\bar{q}$ annihilation, while the remainder originates from gluon-gluon fusion. In 2006, the DØ experiment found evidence for the second Standard Model (SM) mode of production, i.e. the electroweak production of single top quarks. In the SM, a top quark decays $\sim 100\%$ of the time into a W boson and a b-quark. The final state decays of the W boson(s) are used to classify the final state signatures of top quark pair production. If 0, 1, or 2 W bosons decay into a charged lepton and a neutrino, the pair production final states are classified as: all-hadronic, lepton+jets, and dilepton. The branching ratio for a particular final state decreases in this same order.

A precise measurement of the top quark mass is very important for several reasons. First of all, the mass of the top quark is a fundamental parameter of the SM. In addition, the mass of the top quark affects other SM observables via radiative corrections, e.g. the mass of the top quark together with the mass of the W boson can be related to the mass of the yet undiscovered Higgs
boson, thus setting indirect constraints on its value. Figure 1 shows the latest 68% CL contours from a global fit to the electroweak data. The fit yields\(^2\) a prediction for the Higgs boson mass of $76^{+33}_{-24}$ GeV/c\(^2\) and, taking into account the LEP-2 direct limits, a 95% CL one-sided upper limit of 182 GeV/c\(^2\).

The mass of the top quark is also measured to be very large compared to the mass of the other quarks, and of the same order as the vacuum expectation value of the Higgs field, thus suggesting that the top quark might play a special role in the mechanism of electroweak symmetry breaking, and therefore lead to signatures of new physics beyond the SM. The Tevatron experiments are aiming at a precise measurement of the top quark mass. The prediction for the total uncertainty on such a measurement is $\sim 1.5(1.0)$ GeV for an integrated luminosity of 2(8) fb\(^{-1}\).

2 Mass Measuring Techniques

The techniques used by the Tevatron collaborations to extract a measurement of the top quark mass fall into two main categories: template and dynamical methods. The template methods typically extract one value of the top quark mass per event from a kinematic fit to the $t\bar{t}$ hypothesis, then compare the data to a combination of background Monte Carlo templates and signal Monte Carlo templates generated for different top quark mass input values. The dynamical methods weight each event according to the quality of agreement with Standard Model top and background differential cross-sections. The event-by-event weight, or probability, integrates the matrix element for the signal or background processes (LO calculations are used), the parton distributions functions, and the transfer functions. The transfer functions map parton-level variables to reconstructed objects variables. Many variations exist that combine template and dynamical techniques.

As for all other top quark physics measurements, the determination of the mass poses several challenges. Given the complexity of the final state, top quark physics exercises the understanding of all detector components. Top quark production is a rare process, with significant backgrounds. Experimentally, jets are measured, not quarks (with hadronization, radiation, detector effects to be taken into account when trying to calibrate jets and relate jet energies to parton energies). Also, experiments can only infer the presence of a neutrino from the measurement of missing transverse energy. Systematic uncertainties reflect these challenges, and with the increased statistics brought by Run 2 of the Tevatron, the main measurements in top physics are starting to see comparable statistical and systematic uncertainties. With larger and larger data sets
becoming available, the key to a precise measurement of the top quark mass is the control of systematic effects. There are several handles on systematic uncertainties that have been used by the Tevatron collaborations. The uncertainty associated with the determination of the jet energy scale (JES) is predominant in nearly all of the top quark mass measurements. The JES systematic uncertainty can be reduced, in lepton+jets and all-hadronic final states, with the in-situ calibration of the hadronic W mass in top decays. Identification of b−jets with b−tagging can be used to reduce backgrounds as well as combinatorial effects, at the cost of a decrease in statistical power. Most systematic uncertainties are also expected to decrease with the increased statistics of control samples.

3 Results

Within the scope of this contribution it would be impossible to carefully characterize all of the measurement of the top quark mass performed by CDF and DØ since the beginning of Run 2, therefore only the most recent and/or the most precise results in the lepton+jets, dilepton, and all-hadronic final states are reviewed. Most of the results shown here use a data set corresponding to an integrated luminosity of 1 fb⁻¹.

3.1 Top quark mass measurement in the lepton+jets channel

At leading order, the lepton+jets channel final state signature consists of one high \( p_T \) lepton (electron and muon results are considered here), significant \( E_T^{miss} \) from the W decay's neutrino, and four jets, two of which originate from b−quarks (0, 1, 2 b−tags can be required). This channel has good statistics and manageable backgrounds, from W+jets and QCD multijet processes. An in-situ calibration of light quark jets can be performed using the hadronically decaying W mass.

Both Tevatron experiments use a dynamical method in the determination of the most precise value of the top quark mass in this channel. The so-called Matrix Element (ME) method was pioneered by DØ with a re-analysis of the Run 1 data. It makes maximal use of the information in each event by calculating, for each event, a probability of being signal or background based on the matrix elements for the respective processes. All event probabilities are combined in a likelihood which is maximized as a function of three parameters: top quark mass, relative JES, and top signal fraction, \( f_{top} \). The ME measurement requires exactly four jets in the final state.

A new ME measurement by DØ was shown at this conference in the lepton+jets channel. The simultaneous fit to the top quark mass, relative JES, and top fraction yields a value for the top quark mass of \( M_{top} = 170.5 \pm 2.4(\text{stat.} + \text{JES}) \pm 1.2(\text{sys.}) \) GeV/c² with an integrated luminosity of 0.9 fb⁻¹. This measurement requires at least one jet to be b−tagged and it is in good agreement with the result of a similar, but untagged, analysis where the extracted top quark mass is \( M_{top} = 170.5 \pm 2.5(\text{stat.} + \text{JES}) \pm 1.4(\text{sys.}) \) GeV/c². Note that, since the hadronically decaying W's in the decay of top events are used to constraint the JES, the JES error depends on the available statistics and it is incorporated in the statistical error. Figure 2 shows the 2-dimensional likelihood contours as a function of the top quark mass and relative JES for this measurement.

CDF extracts \( a \) a value for the top quark mass, with a similar integrated luminosity and with the same method (ME): \( M_{top} = 170.9 \pm 2.2(\text{stat.} + \text{JES}) \pm 1.4(\text{sys.}) \) GeV/c².

An additional method used by the DØ collaboration is the Ideogram method, which uses the same kinematic fitting and discriminant as basic template analyses, but an event by event likelihood where each event gives a distribution of masses. This method has the advantage of being less CPU intensive than the ME method. A result was recently published with this method, based on a 0.4 fb⁻¹ data sample, yielding a value for the top quark mass of \( M_{top} = 173.7 \pm 4.4(\text{stat.} + \text{JES})^{+1.21}_{-0.35}(\text{sys.}) \) GeV/c².
3.2 Top quark mass measurement in the dilepton channel

The event signature typical of the dilepton channel at LO consists of a pair of high \( p_T \) charged leptons (ee, \( \mu \mu \), or \( e\mu \) are considered here), two jets originating from \( b \)-quarks (\( b \)-tagging can be incorporated) and significant \( E_T^{\text{miss}} \). The backgrounds are low (primarily from diboson and \( W, Z + \text{jets} \) production), even without requiring one jet or more to be \( b \)-tagged. The challenge of this channel is the presence of two neutrinos.

Both template and dynamical methods have been applied to the determination of the top quark mass in the dilepton channel.

Due to the presence of two neutrinos in the event, a kinematic fit of a dilepton event to the \( t\bar{t} \) hypothesis is under-constrained. The template methods assume values for certain variables in order to extract a solution, and assign weights to the different solutions. Different schemes of the weights exist, and the choice of the weight characterizes the measurement.

\( \mathrm{D}\O \) uses two weighting schemes: the matrix weighting and the neutrino weighting (the latter one has been used by CDF as well). The matrix weighting method scans over top quark masses and assigns a weight to the solution, based on the matrix element predictions for the lepton \( p_T \)'s. The neutrino weighting method, scans over several values for the top quark mass and the \( \eta \)'s of the two neutrinos in the event and assigns a weight (as a function of \( m_{\text{top}} \)) to the solution, based on the agreement of the neutrino \( p_T \)'s and the observed \( E_T^{\text{miss}} \).

A maximum likelihood fit of the data to signal and background templates is then performed for all methods in order to extract the top quark mass value and its statistical uncertainty. Since there is no constraint on the JES from a hadronically decaying \( W \), the JES error is part of the systematic uncertainty.

\( \mathrm{D}\O \) performed a new measurement\(^7\) with the neutrino weighting method and an integrated luminosity of 1 fb\(^{-1}\), yielding a result for the top quark mass of: \( M_{\text{top}} = 172.5 \pm 5.8 \text{(stat.)} \pm 5.5 \text{(sys.)} \text{ GeV}/c^2 \). An earlier result\(^8\) with the matrix weighting method on the \( e\mu + \text{jets} \) final state alone yields a consistent result of: \( M_{\text{top}} = 177.7 \pm 8.8 \text{(stat.)} \pm 5.7 \text{(sys.)} \text{ GeV}/c^2 \).

The Matrix Element method is applied to the determination of the top quark mass in the dilepton final state by CDF. The measurement\(^9\) is performed on 1 fb\(^{-1}\) and uses a per-event probability for the mass as a weighted sum of the differential cross section for LO \( t\bar{t} \) production and of the differential cross sections for background processes. A posterior probability density (Figure 3) is formed as the product of a flat prior and the joint event likelihood. The mean and \( \sigma \) of the posterior probability correspond to the value of the top quark mass and its uncertainty.
Figure 3: The joint probability density for the CDF ME measurement of the Top quark mass in the dilepton channel (1030 pb$^{-1}$).

The extracted value of the top quark mass is $M_{\text{top}} = 164.5 \pm 3.9(\text{stat.}) \pm 3.9(\text{sys.})$ GeV/c$^2$, the most precise dilepton measurement. This is in good agreement with a similar measurement which requires at least one jet to be tagged as a $b$-jet: $M_{\text{top}} = 167.3 \pm 4.6(\text{stat.}) \pm 3.8(\text{sys.})$ GeV/c$^2$. The major contributor to the systematic uncertainty in all of the dilepton measurements is the JES uncertainty.

### 3.3 Top quark mass measurement in the all-hadronic channel

The all-hadronic channel has the highest branching ratio and a final state signature, at LO, of six jets, two of which originate from $b$-quarks. Since the QCD multi-jet background in this channel is rather large, at least one jet is required to be $b$-tagged. Selection criteria based on the specific topology of the signal events are also applied to further reduce the background. This channel contains two hadronically decaying W’s in the decay of top and anti-top quarks, thus allowing a similar in-situ calibration of the jet energy scale to the lepton+jets channel.

CDF performed a measurement of the top quark mass in this channel using a Matrix Element assisted template method\textsuperscript{10}. This analysis uses 2-dimensional (in the top quark mass and JES) signal templates derived from ME calculations and background templates modeled on data. Figure 4 shows the extracted top quark mass distribution for single $b$-tagged and for double $b$-tagged events. Based on a 0.9 fb$^{-1}$ data sample, the top quark mass is measured to be: $M_{\text{top}} = 171.1 \pm 3.7(\text{stat.} + \text{JES}) + 2.1(\text{sys.})$ GeV/c$^2$.

Although this channel is the most challenging one, it must be noted that the precision of the measurement is competitive with the precision of the measurements in the other channels.

### 3.4 The top quark mass world average

The measurements of the top quark mass from different channels and the two Tevatron experiments are combined\textsuperscript{11}. The resulting world average for the top quark mass is: $M_{\text{top}} = 170.9 \pm 1.1(\text{stat.}) \pm 1.5(\text{sys.})$. The total uncertainty is 1.8 GeV, the relative uncertainty is 1.1%. The precision of the measurement gives confidence that a $\sim 1$ GeV total uncertainty can be achieved with the full integrated luminosity of Run 2 of the Tevatron. The impact of the new Spring 2007 top quark mass world average on the indirect Higgs boson mass constraints was discussed in Section 1.
Figure 4: Results of the unbinned likelihood fit for the top quark mass in the all-hadronic channel at CDF, requiring one $b$-tag (left) or two $b$-tags (right) in the event.

4 Conclusions and outlook

Results on the measurement of the top quark mass at the Tevatron were presented for datasets corresponding to integrated luminosities up to 1 fb$^{-1}$. All of the 1 fb$^{-1}$ measurements are currently converging and analyses of the 2 fb$^{-1}$ dataset have started. Measurements now extend to final states which were once considered challenging, such as the all-hadronic mode, with results competitive in precision with other channels. The current relative uncertainty on the combined value of the top quark mass from the Tevatron is 1.1%. The Tevatron collaborations are aiming at a combined $\sim$1 GeV total uncertainty ($<1\%$ in relative uncertainty) by the end of Run 2 of the Tevatron. At this level of precision, a discussion about theoretical uncertainties in how to interpret this measurement will be needed. The current excellent performance of the Tevatron and of the CDF and DØ detectors are the key to precise measurements in top physics.

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SM SCALAR AND EXTRA SINGLET(S)

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I discuss the question whether it is possible that the LHC will find no signal for the Higgs particle. It is argued that in this case singlet scalars should be present that could play an important role in astroparticle physics. A critical view at the existing electroweak data shows that this possibility might be favored over the simplest standard model. In this case one needs the ILC in order to study the Higgs sector.

1 Introduction

The standard model gives a good description of the bulk of the electroweak data. Only a sign of the Higgs particle is missing at the moment. The Higgs field is necessary in order to make the theory renormalizable, so that predictions are possible and one can really speak of a theory. A complete absence of the Higgs field would make the theory non-renormalizable, implying the existence of new strong interactions at the TeV scale. Therefore one is naively led to the so-called no-lose theorem. This theorem says that when one builds a large energy hadron collider, formerly the SSC now the LHC, one will find new physics, either the Higgs particle or otherwise new strong interactions. Since historically no-theorems have a bad record in physics one is naturally tempted to try to evade this theorem. So in the following I will try to find ways by which the LHC can avoid seeing any sign of new physics.

At the time of the introduction of the no-lose theorem very little was known about the Higgs particle. Since then there have been experiments at LEP, SLAC and the Tevatron, that give information on the Higgs mass. Through precise measurements of the W-boson mass and various asymmetries one can get constraints on the Higgs mass. The Higgs mass enters into the prediction of these quantities via radiative corrections containing a virtual Higgs exchange. Moreover at LEP-200 the direct search gives a lower limit of 114.4 GeV. The situation regarding the precision tests is not fully satisfactory. The reason is that the Higgs mass implied by
the forward-backward asymmetry $A_{FB}(b)$ from the bottom quarks is far away from the mass implied by the other measurements, that agree very well with each other. No model of new physics appears to be able to explain the difference. From $A_{FB}(b)$ one finds $m_H = 488^{+426}_{-219}$ GeV with a 95% lower bound of $m_H = 181$ GeV. Combining the other experiments one finds $m_H = 51^{+12}_{-11}$ GeV with a 95% upper bound of $m_H = 109$ GeV. The $\chi^2$ of the latter fit is essentially zero. Combining all measurements gives a bad fit. One therefore has a dilemma. Keeping all data one has a bad fit. Ignoring the $b$-data the standard model is ruled out. In the last case one is largely forced towards the extended models that appear in the following. Accepting a bad fit one has somewhat more leeway, but the extended models are still a distinct possibility.

2 Is a very heavy Higgs boson possible?

One way to avoid seeing the Higgs boson would be if it is too heavy to be produced at the LHC. At first sight this possibility appears to be absurd given the precision data. Even if one takes all data into account there is an upper limit of $m_H = 190$ GeV. However the question is surprisingly difficult to answer in detail. The reason is that the Higgs mass is not a free parameter in the Lagrangian. Because of the spontaneous symmetry breaking the Higgs mass is determined by its self-coupling $\lambda$ and the vacuum expectation value $f$: $m_H^2 = \lambda f^2$. This means that a heavy Higgs boson is strongly interacting. Therefore higher-loop effects can become important. These effects give corrections to the precision measurements with a behaviour $m_H^2(\text{loop})^{-1}$. These effects can in principle cancel the one-loop $\log(m_H)$ corrections, on which the limits are based. Therefore one could have the following situation: the strong interactions compensate for the loop effects, so that from the precision measurements the Higgs appears to have a mass of 50 GeV. At the same time the Higgs is so heavy that one does not see it at the LHC. For this to happen the Higgs mass would have to be about 3 TeV. Detailed two-loop $\int,\int,\int,\int$ and non-perturbative $1/N$ calculations $\int,\int$ have shown that the first important effects are expected at the three-loop level. The important quantity is the sign of the three-loop correction compared to the one-loop correction. This question was settled in a large calculation that involved the order of half a million Feynman diagrams $\int,\int$. The conclusion is that the strong interactions enhance the effects of a heavy Higgs boson. This conclusion is confirmed by somewhat qualitative non-perturbative estimates $\int,\int$. Therefore the Higgs boson cannot be too heavy to be seen at the LHC.

3 Singlet scalars

3.1 Introduction

If the Higgs boson is not too heavy to be seen the next try to make it invisible at the LHC is to let it decay into particles that cannot be detected. For this a slight extension of the standard model is needed. In order not to effect the otherwise good description of the electroweak data by the standard model one introduces singlet scalars. The presence of singlets will not affect present electroweak phenomenology in a significant way, since their effects in precision tests appear first at the two-loop level and are too small to be seen $\int$. These singlet scalars will not couple to ordinary matter in a direct way, but only to the Higgs sector. It is actually quite natural to expect singlet scalars to be present in nature. After all we know there also exist singlet fermions, namely the right handed neutrino's. The introduction of singlet scalars affects the phenomenology of the Higgs boson in two ways. On the one hand one creates the possibility for the Higgs boson to decay into said singlets, on the other hand there is the possibility of singlet-doublet mixing, which will lead to the presence of more Higgs bosons however with reduced couplings to ordinary matter. In the precision tests this only leads to the replacement of the single Higgs mass by a weighted Higgs mass and one cannot tell the difference between the two cases. Mixing and
invisible decay can appear simultaneously. For didactical purpose I show in the following simple models consisting of pure invisible decay or pure mixing. For a mini-review of the general class of models see ref.\(^7\).

3.2 Invisible decay

When singlet scalars are present it is possible that the Higgs boson decays into these scalars if they are light enough. Such an invisible decay is rather natural, when one introduces the Higgs singlets \(S_i\) as multiplets of a symmetry group\(^\ldots\), for instance \(O(N)\). When the \(O(N)\) symmetry group stays unbroken this leads to an invisibly decaying Higgs boson through the interaction \(\Phi^\dagger \Phi S_i S_i\), after spontaneous breaking of the standard model gauge symmetry. When the \(O(N)\) symmetry stays unbroken the singlets \(S_i\) are stable and are suitable candidates for the dark matter in the universe\(^\ldots\).

To be more concrete let us discuss the Lagrangian of the model, containing the standard model Higgs boson plus an \(O(N)\)-symmetric sigma model. The Lagrangian density is the following:

\[
L_{\text{Scalar}} = L_{\text{Higgs}} + L_S + L_{\text{Interaction}}
\]

\[
L_{\text{Higgs}} = -\frac{1}{2} D_\mu \Phi^\dagger D_\mu \Phi - \frac{\lambda}{8} (\Phi^\dagger \Phi - f^2)^2
\]

\[
L_S = -\frac{1}{2} \partial_\mu S_i \partial^\mu S_i - \frac{1}{2} m_S^2 S_i^2 - \frac{\lambda S}{8 N} (S^2)^2
\]

\[
L_{\text{Interaction}} = -\frac{\omega}{4 N} \bar{S} \Phi^\dagger \Phi
\]

The field \(\Phi = (\sigma + f + i \pi_1, \pi_2 + i \pi_3)\) is the complex Higgs doublet of the standard model with the vacuum expectation value \(\langle 0 | \Phi | 0 \rangle = (f, 0)\), \(f = 246\) GeV. Here, \(\sigma\) is the physical Higgs boson and \(\pi_{1,2,3}\) are the three Goldstone bosons. \(S = (S_1, \ldots, S_N)\) is a real vector with \(\langle 0 | S | 0 \rangle = \bar{S}\). We consider the case, where the \(O(N)\) symmetry stays unbroken, because we want to concentrate on the effects of a finite width of the Higgs particle. Breaking the \(O(N)\) symmetry would lead to more than one Higgs particle, through mixing. After the spontaneous breaking of the standard model gauge symmetry the \(\pi\) fields become the longitudinal polarizations of the vector bosons. In the unitary gauge one can simply put them to zero. One is then left with an additional interaction in the Lagrangian of the form:

\[
L_{\text{Interaction}} = -\frac{\omega f}{2 \sqrt{N}} S^2 \sigma
\]

This interaction leads to a decay into the \(S\) particles, that do not couple to other fields of the standard model Lagrangian. On has therefore an invisible width:

\[
\Gamma_{\text{Higgs}}(\text{invisible}) = \frac{\omega^2}{32 \pi} \frac{f^2}{m_{\text{Higgs}}} (1 - 4 m_S^2 / m_{\text{Higgs}}^2)^{1/2}
\]

This width is larger than the standard model width even for moderate values of \(\omega\), because the standard model width is strongly suppressed by the Yukawa couplings of the fermions. Therefore the Higgs boson decays predominantly invisibly with a branching ratio approximating 100%. Moreover one cannot exclude a large value of \(\omega\). In this case the Higgs is wide and decaying invisibly. This explains the name stealth model for this kind of Higgs sector.

However, is this Higgs boson undetectable at the LHC? Its production mechanisms are exactly the same as the standard model ones, only its decay is in undetectable particles. One therefore has to study associated production with an extra Z-boson or one must consider the
vector-boson fusion channel with jet-tagging. Assuming the invisible branching ratio to be large and assuming the Higgs boson not to be heavy, as indicated by the precision tests, one still finds a significant signal. Of course one cannot study this Higgs boson in great detail at the LHC. For this the ILC would be needed, where precise measurements are possible in the channel $e^+e^- \rightarrow ZH$.

### 3.3 Mixing: fractional Higgses

Somewhat surprisingly it is possible to have a model that has basically only singlet-doublet mixing even if all the scalars are light. If one starts with an interaction of the form $H \Phi \Phi$, where $H$ is the new singlet Higgs field and $\Phi$ the standard model Higgs field, no interaction of the form $H^3$, $H^4$ or $H^2 \Phi \Phi$ is generated with an infinite coefficient. At the same time the scalar potential stays bounded from below. This means that one can indeed leave these dimension four interactions out of the Lagrangian without violating renormalizability. This is similar to the non-renormalization theorem in supersymmetry that says that the superpotential does not get renormalized. However in general it only works with singlet extensions. As far as the counting of parameters is concerned this is the most minimal extension of the standard model, having only two extra parameters.

The simplest model is the Hill model:

$$L = -\frac{1}{2}(D_\mu \Phi^\dagger)(D_\mu \Phi) - \frac{1}{2}(\partial_\mu H)^2 - \frac{\lambda_0}{8}(\Phi^\dagger \Phi - f_0^2)^2 - \frac{\lambda_1}{8}(2f_1H - \Phi^\dagger \Phi)^2$$  \hspace{1cm} (7)

Working in the unitary gauge one writes $\Phi^\dagger = \sigma 0$, where the $\sigma$-field is the physical standard model Higgs field. Both the standard model Higgs field $\sigma$ and the Hill field $H$ receive vacuum expectation values and one ends up with a two-by-two mass matrix to diagonalize, thereby ending with two masses $m_-$ and $m_+$ and a mixing angle $\alpha$. There are two equivalent ways to describe this situation. One is to say that one has two Higgs fields with reduced couplings $g$ to standard model particles:

$$g_- = g_{SM} \cos(\alpha), \quad g_+ = g_{SM} \sin(\alpha)$$  \hspace{1cm} (8)

Because these two particles have the quantum numbers of the Higgs particle, but only reduced couplings to standard model particles one can call them fractional Higgs particles. The other description, which has some practical advantages is not to diagonalize the propagator, but simply keep the $\sigma - \sigma$ propagator explicitly. One can ignore the $H - \sigma$ and $H - H$ propagators, since the $H$ field does not couple to ordinary matter. One simply replaces in all experimental cross section calculations the standard model Higgs propagator by:

$$D_{\sigma\sigma}(k^2) = \frac{\cos^2(\alpha)}{(k^2 + m_-^2)} + \frac{\sin^2(\alpha)}{(k^2 + m_+^2)}$$  \hspace{1cm} (9)

The generalization to an arbitrary set of fields $H_k$ is straightforward, one simply replaces the singlet-doublet interaction term by:

$$L_{\Phi \Phi} = -\sum_k \frac{\lambda_k}{8}(2f_k H_k - \Phi^\dagger \Phi)^2$$  \hspace{1cm} (10)

This will lead to a number of (fractional) Higgs bosons $H_i$ with reduced couplings $g_i$ to the standard model particles such that

$$\sum_i g_i^2 = g_{SM}^2$$  \hspace{1cm} (11)
3.4 A higher dimensional Higgs boson

The mechanism described above can be generalized to an infinite number of Higgses. The physical Higgs propagator is then given by an infinite number of very small Higgs peaks, that cannot be resolved by the detector. Ultimately one can take a continuum limit, so as to produce an arbitrary line shape for the Higgs boson, satisfying the Källén-Lehmann representation.

\[ D_{\sigma\sigma}(k^2) = \int ds \rho(s)/(k^2 + \rho(s) - i\epsilon) \]  

(12)

One has the sum rule: \( \int \rho(s) ds = 1 \), while otherwise the theory is not renormalizable and would lead to infinite effects for instance on the LEP precision variables. Moreover, combining mixing with invisible decay, one can vary the invisible decay branching ratio as a function of the invariant mass inside the Higgs propagator. There is then no Higgs peak to be found any more. The general Higgs propagator for the Higgs boson in the presence of singlet fields is therefore determined by two function, the Källén-Lehmann spectral density and the s-dependent invisible branching ratio. Unchanged compared to the standard model are the relative branching ratio’s to standard model particles.

Given the fact that the search for the Higgs boson in the low mass range heavily depends on the presence of a sharp mass peak, this is a promising way to hide the Higgs boson at the LHC. However the general case is rather arbitrary and unelegant and ultimately involves an infinite number of coupling constants. The question is therefore whether there is a more esthetic way to generate such a spread-out Higgs signal, without the need of a large number of parameters. Actually this is possible. Because the \( H\Phi^4\Phi \) interaction is superrenormalizable one can let the \( H \) field move in more dimensions than four, without violating renormalizability. One can go up to six dimensions. The precise form of the propagator will in general depend on the size and shape of the higher dimensions. The exact formulas can be quite complicated. However it is possible that these higher dimensions are simply open and flat. In this case one finds simple formulas. One has for the generic case a propagator of the form:

\[ D_{\sigma\sigma}(q^2) = \left[ q^2 + M^2 - \frac{\mu_{\phi}^2}{\mu_{\phi}^2} \left( q^2 + m^2 \right) \right]^{-1}. \]  

(13)

For six dimensions one needs a limiting procedure and finds:

\[ D_{\sigma\sigma}(q^2) = \left[ q^2 + M^2 + \frac{\mu_{\phi}^2}{\mu_{\phi}^2} \log \left( \frac{q^2 + m^2}{\mu_{\phi}^2} \right) \right]^{-1}. \]  

(14)

The parameter \( M \) is a four-dimensional mass, \( m \) a higher-dimensional mass and \( \mu_{\phi} \) a higher-to-lower dimensional mixing mass scale. When one calculates the corresponding Källén-Lehmann spectral densities one finds a low mass peak and a continuum that starts a bit higher in the mass. The location of the peak is given by the zero of the inverse propagator. Because this peak should not be a tachyon, there is a constraint on \( M, m, \mu_{\phi} \), that can be interpreted as the condition that there is a stable vacuum.

Explicitly one finds for \( d = 5 \) the Källén-Lehmann spectral density:

\[ \rho(s) = \theta(m^2 - s) \frac{2(m^2 - s_{\text{peak}})^{3/2}}{2(m^2 - s_{\text{peak}})^{3/2} + \mu_{\phi}^2} \delta(s - s_{\text{peak}}) \]

\[ + \quad \frac{\mu_{\phi}^2}{s - m^2} \frac{\mu_{\phi}^2}{(s - m^2)(s - M^2)^2} \delta(s - s_{\text{peak}}) \]  

(15)

For \( d = 6 \) one finds:

\[ \rho(s) = \theta(m^2 - s) \frac{m^2 - s_{\text{peak}}}{m^2 + \mu_{\phi}^2 - s_{\text{peak}}} \delta(s - s_{\text{peak}}) \]

\[ + \quad \theta(s - m^2) \frac{\mu_{\phi}^2}{s - M^2 - \mu_{\phi}^2} \frac{\mu_{\phi}^2}{\mu_{\phi}^2 \log((s - m^2)/\mu_{\phi}^2)^{1/2} + \mu_{\phi}^2} \delta(s - s_{\text{peak}}) \]  

(16)
If one does not introduce further fields no invisible decay is present. If the delta peak is small enough it will be too insignificant for the LHC search. The continuum is in any case difficult to see. There might possibly be a few sigma signal in the $\tau$-sector. However if one adds to this model some scalars to account for the dark matter, this will water down any remnant signal to insignificance.

4 Comparison with the LEP-200 data

We now confront the higher dimensional models with the results from the direct Higgs search at LEP-200. Within the pure standard model the absence of a clear signal has led to a lower limit on the Higgs boson mass of 114.4 GeV at the 95% confidence level. Although no clear signal was found the data have some intriguing features, that can be interpreted as evidence for Higgs bosons beyond the standard model. There is a $2.3 \sigma$ effect seen by all experiments at around 98 GeV. A somewhat less significant $1.7 \sigma$ excess is seen around 115 GeV. Finally over the whole range $s^{1/2} > 100$ GeV the confidence level is less than expected from background. We will interpret these features as evidence for a spread-out Higgs-boson. The peak at 98 GeV will be taken to correspond to the delta peak in the Källén-Lehmann density. The other excess data will be taken as part of the continuum, that will peak around 115 GeV.

We start with the case $d = 5$. The delta-peak will be assumed to correspond to the peak at 98 GeV, with a fixed value of $g_{58}^2$. Ultimately we will vary the location of the peak between $95 \text{ GeV} < m_{\text{peak}} < 110 \text{ GeV}$ and $0.056 < g_{58}^2 < 0.144$. After fixing $g_{58}^2$ and $m_{\text{peak}}$ we have one free variable, which we take to be $\mu_{\text{h}d}$. If we also take a fixed value for $\mu_{\text{h}d}$ all parameters and thereby the spectral density is known. We can then numerically integrate the spectral density over selected ranges of $s$. The allowed range of $\mu_{\text{h}d}$ is subsequently determined by the data at 115 GeV. Since the peak at 115 GeV is not very well constrained, we demand here only that the integrated spectral density from $s_{\text{down}} = (110 \text{ GeV})^2$ to $s_{\text{up}} = (120 \text{ GeV})^2$ is larger than 30%. This condition, together with formula (15), which implies:

$$\rho(s) < \frac{(s - m^2)^{1/2}}{\pi \mu_{\text{h}d}^2},$$

leads to the important analytical result:

$$\frac{2}{3\pi \mu_{\text{h}d}^2} | (s_{\text{up}} - m_{\text{peak}}^2)^{3/2} - (s_{\text{down}} - m_{\text{peak}}^2)^{3/2} | > 0.3$$

This implies $\mu_{\text{h}d} < 53 \text{ GeV}$. Using the constraint from the strength of the delta-peak, it follows that the continuum starts very close to the peak, the difference being less than 2.5 GeV. This allows for a natural explanation, why the CL for the fit in the whole range from 100 GeV to 110 GeV is somewhat less than what is expected by pure background. The enhancement can be due to a slight, spread-out Higgs signal. Actually when fitting the data with the above conditions, one finds for small values of $\mu_{\text{h}d}$, that the integrated spectral density in the range 100 GeV to 110 GeV can become rather large, which would lead to problems with the 95% CL limits in this range. We therefore additionally demand that the integrated spectral density in this range is less than 30%. There is no problem fitting the data with these conditions. As allowed ranges we find:

$$95 \text{ GeV} < m < 101 \text{ GeV}$$
$$111 \text{ GeV} < M < 121 \text{ GeV}$$
$$26 \text{ GeV} < \mu_{\text{h}d} < 49 \text{ GeV}$$

(19)
We now repeat the analysis for the case $d = 6$. The analytic argument gives the result:

\[
\frac{\sigma_{\text{up}} - \sigma_{\text{down}}}{\pi \mu_{\text{h}}^2} > 0.3 \tag{20}
\]

which implies $\mu_{\text{h}} < 28 \text{ GeV}$. Because of this low value of $\mu_{\text{h}}$, it is difficult to get enough spectral weight around 115 GeV and one also tends to get too much density below 110 GeV. As a consequence the fit was only possible in a restricted range. Though not quite ruled out, the six-dimensional case therefore seems to be somewhat disfavoured compared to the five-dimensional case. As a consequence the fit was only possible in a restricted range. We found the following limits:

\begin{align*}
95 \text{ GeV} &< m < 101 \text{ GeV} \\
106 \text{ GeV} &< M < 111 \text{ GeV} \\
22 \text{ GeV} &< \mu_{\text{h}} < 27 \text{ GeV} \tag{21}
\end{align*}

5 Conclusion

We are now in a position to answer the following questions. Is it possible to have a simple model that:

a) Is consistent with the precision data, even with the strong condition $m_H < 109 \text{ GeV}$?

b) explains the LEP-200 Higgs search data?

c) has a dark matter candidate?

d) gives no Higgs signal at the LHC?

Given the above discussion, the answer is clearly yes, which leads to the question whether such a model is likely to be true. This is rather difficult to answer decisively. It depends on how significant the evidence in the data is, in particular in the LEP-200 Higgs search data. This significance is hard to estimate, since the data were not analyzed with this type of model in mind. Taking the situation at face value the spread-out singlet models appear to be the only way to satisfy the experimental constraints. In that case one is led to the conclusion that the LHC will not see a signal for the Higgs boson.

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Higgs Searches at the Tevatron

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Recent preliminary results obtained by the CDF and DØ Collaborations on searches for the Higgs boson in $p\bar{p}$ collisions at $\sqrt{s} = 1.96$ TeV at the Fermilab Tevatron collider are discussed. The data, corresponding to integrated luminosities of about 1 fb$^{-1}$, show no excess of a signal above the expected background in any of the decay channels examined. Instead, upper limits at 95% Confidence Level on the cross section are established. Further, a combined Standard Model Tevatron cross section limit is presented.

1 Introduction

The Higgs boson is the last missing particle in the Standard Model (SM). Its mass is not determined by the SM, there are however several experimental constraints which bound the Higgs mass to values which are within the reach of the Tevatron collider. Lower bounds are given from direct searches at LEP2. These results exclude Higgs masses below 114.4 GeV at the 95% Confidence Level (C.L.)$^1$. An upper bound on the Higgs mass is obtained by global electroweak fits. Especially radiative corrections to the $W$ mass from the Higgs and top quark play an important role. New precision measurements of the $W$ mass$^2$ and the top mass$^3$ from the Tevatron favor a light SM Higgs boson and yield an upper value of 144 GeV at 95% C.L. (or 182 GeV if the LEP2 limit is included)$^4$.

2 Experimental environment

The Higgs searches are crucially dependent on performance of the Tevatron accelerator and detectors. Both, CDF and DØ detectors are currently performing close to their optimal design values, taking data with an efficiency of about 90%. The present Tevatron performance is matching the design values in terms of the current weekly integrated and peak luminosity. As

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$^a$Presented at Rencontres de Moriond EW 2007, 10-17 March 2007, La Thuile.
of today, more than 2.5 fb$^{-1}$ have been delivered, with weekly integrated luminosity routinely reaching 50 pb$^{-1}$. If the accelerator keeps following the designed luminosity evolution, an integrated luminosity of about 8 fb$^{-1}$ will be achieved by the end of 2009, increasing the potential for a Higgs discovery at the Tevatron significantly.

3 Standard Model Higgs searches

Production cross sections for the SM Higgs boson at the Tevatron are rather small. They depend on the Higgs mass and are about 0.1 – 1 pb in the mass range of 100 – 200 GeV. The largest production cross section comes from gluon fusion, where the Higgs is produced via a quark loop. The second largest cross section, almost an order of magnitude smaller, is the associated production with vector bosons. At the mass range covered by the Tevatron, below 135 GeV the highest branching ratio is given by the decay to $b\bar{b}$ pairs and for masses above 135 GeV the Higgs boson decays mainly to $WW$ pairs.

These production and decay properties lead to the following search strategy at the Tevatron:

- For masses below 135 GeV the main search channels are the associated productions with vector bosons where the Higgs decays into $b\bar{b}$ pairs. In order to isolate the main background processes to these channels, an efficient b-tagging algorithm and a good dijet mass resolution are essential. The same final state produced via the gluon fusion process leads to a higher cross section but is overwhelmed by the huge multijet QCD background at a hadron collider.

- For masses above 135 GeV the search is mainly focused on the gluon fusion production process where the Higgs decays into $WW$ pairs.

3.1 $WH \rightarrow \ell
\nu b\bar{b}$, $\ell = e, \mu$

For SM Higgs searches the most sensitive production channel at the Tevatron for a Higgs mass below 135 GeV is the associated production of a Higgs boson with a $W$ boson. Dominant backgrounds to the $WH$ signal are $W +$ heavy flavor production, $t\bar{t}$ and single-top quark production. Both, CDF and DØ performed cut based analyses with a rather similar approach. Both, electron and muon channels are studied here. The channels are separated in events having exactly one "tight" b-tagged jet (ST), and those having two "loose" b-tagged jets (DT) (with no overlap). The resulting four channels are analyzed independently to optimize the sensitivity and are later combined. Both experiments select events with isolated electrons or muons with $p_T > 20$ GeV, require missing transverse energy above 20 GeV and two jets with $p_T > 20$ GeV (DO) or $p_T > 15$ GeV (CDF). Cross section limits are derived from the invariant dijet mass distribution of the four individual analyses of each experiment and later combined. For $m_H = 115$ GeV the observed (expected) limit is 1.3 (1.1) pb at DØ$^S$ and 3.4 (2.2) pb at CDF$^S$, to be compared to the Standard Model cross section expectation of 0.13 pb. Thus the best expected measurement is a factor 8.8 higher than the SM expectation.

DØ analyzed this channel also with the Matrix Element technique to separate signal from background. Like in the cut based analysis the four channels ($e, \mu, ST, DT$) are analyzed separately and later combined. The matrix-element-based technique attempts to make use of all the available kinematic information in the event to separate signal and background. Therefore leading order Matrix Elements are used to compute the event probabilities for signal and background. The present selection criteria is based on the single top search and will be optimized in the future. Although this selection is not optimal for $WH$, the sensitivity of this search is similar to the sensitivity of the cut-based analysis and will improve with an optimized selection. For $m_H = 115$ GeV the observed (expected) limit is 1.7 (1.2) pb with this present approach$^7$. Limits for other Higgs masses together with the cut-based results are displayed in Fig.1.
3.2 $ZH \to ℓb\bar{b}$, $ℓ = e, μ$

Similarly to $WH$ the Higgs boson can be produced associated with the $Z$ boson. First we focus on the channel where the $Z$ boson decays to a pair of electrons or muons with opposite sign. Here the $Z$ boson is reconstructed and identified from a pair of high $p_T$ leptons with an invariant mass constraint. Events are required to have $b$-tagged jets. The dominant backgrounds result from the associated production of a $Z$ boson with jets, among which the $Zb\bar{b}$ production is an irreducible background. Other main backgrounds are $tt$, $WZ$, $ZZ$, and multijet production from QCD processes.

In the search at $DΩ$ at least two $b$-tagged jets are required. Cross section limits are then derived from the dijet invariant mass distribution within a search window. At CDF only 1 $b$-tagged jet is required. After this, a two dimensional Neural Network discriminates against the two largest backgrounds which are $Z +$ jets and $tt$. Limits are derived from the Neural Network distribution. For $m_H = 115$ GeV the observed (expected) limit is $2.7 (2.8)$ pb at $DΩ^8$ and $2.2 (1.9)$ at CDF$^9$, to be compared to the Standard Model cross section expectation of 0.08 pb.

3.3 $ZH \to ννb\bar{b}$, $WH \to (ℓ^±)νb\bar{b}$

The $ZH \to ννb\bar{b}$ channel benefits from the large $Z \to νν$ branching ratio. However it is challenging at hadron colliders due to the absence of visible leptons and the presence of only two jets in the final state. The two $b$-jets from the Higgs are boosted along the direction of the Higgs momentum and so tend to be more acoplanar than the dijet background. There are two major sources of background: physics backgrounds such as $Z$+jets, $W$+jets, electroweak diboson production or top quark production with missed leptons and jets and the instrumental background resulting from calorimeter mismeasurements which can lead to high $E_T$ signals with the presence of jets from QCD processes.

A result on this search channel was presented from CDF. Selecting events with a large $E_T > 75$ GeV and high $p_T$ $b$-tagged jets (leading jet $p_T > 60$ GeV), vetoing events with isolated leptons or where the missing $E_T$ is aligned in φ with jets eliminate much of the physics background. Two separate analyses are optimized for one or two $b$-tagged samples and later combined. Since the $WH$ channel with an undetected lepton has the same signature these events are taken into account in this search channel. For $m_H = 115$ GeV the expected limit at CDF is a factor 15 higher than the Standard Model expectation$^{10}$. 

3.4 $H \to WW^{(*)} \to ℓ^+ℓ^-ν\bar{ν}$, $ℓ = e, μ$

At Higgs masses above 135 GeV the biggest branching ratio is the decay to $WW$ pairs. With only leptons and missing energy in the final state the main background is $WW$ production without a large overlapping QCD background. Both, CDF and $DΩ$ analyzed this channel for the three combinations of electron and muon final states. Later the cross section limits have been combined.

The search strategy is to look for two high $p_T$, isolated, opposite sign leptons, require large missing transverse energy and veto on events with jets. Finally, the spin correlations in the decay of the Higgs boson are used. The leptons of the Higgs decay tend to have a small opening angle, whereas leptons from most of the backgrounds are expected to be back-to-back. Thus a cut on the opening angle between the leptons in the transverse plane $Δφ_{ll}$ is mainly used to discriminate against the dominant $WW$ background. Since the Higgs mass cannot be directly reconstructed due to the neutrinos in the final state, the cross section limit is derived from the $Δφ_{ll}$ distribution. For $m_H = 160$ GeV, which yields the best sensitivity, the expected limit at $CDF^{11}$ is a factor 6 and at $DΩ^{12,13}$ a factor of 5 higher than the Standard Model expectation.
3.5 Combined Standard Model Higgs limits

The above presented channels can be combined which leads to a much more sensitive cross section limit throughout the whole discussed mass range. Both, DØ and CDF released results on the SM Higgs combination, the obtained results can be found in [14]. A further, important increase of the sensitivity can be gained from a combination of the CDF and DØ results. Such a first Tevatron combination limit was released Summer 2006, the result is plotted in Fig.2. The expected combined limits are a factor of 7.5 at $m_H = 115$ GeV and a factor of 4 at $m_H = 160$ GeV away from the Standard Model expected cross sections. It should be stressed that this result does not include CDF’s new 1 fb$^{-1}$ high mass results and it does not include any of DØ’s new 1 fb$^{-1}$ low mass results yet. Further significant improvements are expected when all the 1 fb$^{-1}$ results will be included. Such a new Tevatron combination is planned for the Summer 2007.

4 MSSM Higgs searches

The Minimal Supersymmetric Standard Model (MSSM) predicts two Higgs doublets leading to five Higgs bosons: a pair of charged Higgs boson ($H^\pm$); two neutral CP-even Higgs bosons ($h, H$) and a CP-odd Higgs boson ($A$). At tree level, the Higgs sector of the MSSM is fully described by two parameters, which are chosen to be the mass of the CP-odd Higgs, $m_A$, and tan $\beta$, the ratio of the vacuum expectation values of the two Higgs doublets. The Higgs production cross-section is enhanced in the region of low $m_A$ and high tan $\beta$ due to the enhanced Higgs coupling to down-type fermions. This makes it possible to search in the MSSM for $\tau\tau$ final states, which would be very challenging in the SM due to the large irreducible background of $Z \rightarrow \tau\tau$. In the low $m_A$, high tan $\beta$ region of the parameter space, Tevatron searches can therefore probe several MSSM benchmark scenarios extending the search regions covered by LEP [15].

Both, CDF and DØ performed a search for the neutral MSSM Higgs decaying to $\tau$ pairs, where one of the $\tau$-leptons is decaying in the leptonic and the other one in the hadronic mode. DØ’s result covers so far only the $\mu$-channel, CDF’s result is a combination of the electron and
muon channels, including $\tau_e\tau_\mu$.

A set of Neural Networks (NN) is used at DØ to discriminate $\tau$-leptons from jets. An isolated muon is required, separated from the hadronic $\tau$ with opposite sign. A cut on the visible $W$ mass removes most of the remaining $W$ boson background. Further optimized NNs are used for signal discrimination. In the cross section limit calculation the output of the NNs for different $\tau$ types is used.

CDF uses a variable cone size algorithm for $\tau$ discrimination. An isolated muon or electron is required, separated from the hadronic $\tau$ with opposite sign. Most of the $W$ background is removed by a requirement on the relative directions of the visible $\tau$ decay products and the missing transverse energy. Cross section limits are derived from the visible mass distribution.

For both experiments the data is consistent with the background only observation. Exclusion regions in the tan $\beta - m_A$ plane can be derived for different MSSM benchmark scenarios. Both experiments obtained similar results $^{16,17,18}$. In the region of $90 < m_A < 200$ GeV, tan $\beta$ values larger than 40-60 are excluded for the no-mixing and the $m_h^{max}$ benchmark scenarios. Examples of such exclusion regions are shown in Fig.3. In CDF's result the observed limits are weaker than the expectations due to some excess of events in the data sample with a significance of approximately $2\sigma$.

5 Perspectives

Today some single channels have cross section limits similar to the combined Tevatron results obtained half a year ago. With Tevatron's excellent performance matching the designed delivered weekly luminosities, a significant amount of sensitivity will be gained by an increase of the luminosity by about a factor of 8. There is already 2.5 times more data on tape than used for the presented results. In addition, the inclusion of more channels in the Higgs search (for example $\tau$-final states) will gain additional sensitivity. Dijet mass resolution, b-tagging and simulation are important ingredients for Higgs searches and both experiments are continuously improving at these scopes. Still a lot of improvements are expected in analyses techniques. Especially the use of multivariate techniques, like Neural Networks, Decision Trees and Matrix Element analyses.
shall bring further important improvements. DØ’s recent evidence for Single Top production and CDF’s $WZ$ observation is an important milestone in the use of these techniques to discriminate very low rate signals in the presence of substantial backgrounds.

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Electroweak Interactions and Unified Theories

II - Beyond the Standard Model (BSM) and Searches
SEARCHES FOR NEW PHENOMENA WITH LEPTON FINAL STATES AT THE TEVATRON

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Numerous searches for new phenomena have been carried out using data from proton-antiproton collisions at Fermilab's Tevatron. Final states with leptons give signatures which are relatively unique and generally have small backgrounds. We present many of the latest results from the CDF and D0 collaborations from 0.4-1.2 fb$^{-1}$ of data. Topics include supersymmetry, extra gauge bosons, Randall-Sundrum gravitons, excited electrons and neutral, long-lived particles.

New phenomena searches with leptons favor analyses with relatively small or well-understood backgrounds. At the Tevatron, the CDF and D0 collaborations have strong programs of searches with leptonic final states in a wide range of topics. Some of the most recent results are discussed here. All limits presented are at the 95% CL.

1 Charginos and Neutralinos in Trileptons

Supersymmetry has been a popular extension to the standard model for several decades and numerous searches for evidence of any of the superpartners have been carried out. At hadron colliders, a unique signature for supersymmetry comes in the trilepton ("three lepton") final states. If the masses lie in the correct region, proton-antiproton collisions can produce charginos and neutralinos in association:

$$pp \rightarrow \chi^{\pm}_1 \chi^0_2$$  \hspace{1cm} (1)

with decay modes
Figure 1: Limits on SUSY production of charginos and neutralinos. (a) CDF limit using a model of MSSM without slepton mixing and \( m_0 = 60 \text{ GeV} \); (b) DØ limit using three mSUGRA-inspired models.

\[
\tilde{\chi}_1^\pm \rightarrow \ell^\pm \nu \chi_1^0 \quad \tilde{\chi}_2^0 \rightarrow \ell^\pm \ell^\mp \chi_1^0
\]  

(2)

In the standard model, trilepton final states are only produced by rare processes (such as di-bosons) which means these searches will naturally have small backgrounds. The challenge lies in the inefficiency to uniquely identify all three leptons. The solution is to use three search techniques: (1) observe all three leptons; (2) observe two leptons and a third isolated track; (3) observe two same-signed leptons. By combining all three search methods and combinations of electrons, muons and isolated tracks, the experiments improve the sensitivity to discovery.

CDF has performed searches in 14 different channels \(^1\) ranging in luminosity 0.7-1.1 fb\(^{-1}\). While a slight excess of data vs. background is observed, there is no strong evidence of supersymmetry. Therefore limits on the production cross-section times branching ratio can be set. CDF interprets this within three different mSUGRA inspired models with \( m_0 = 60 \text{ GeV} \): (A) mSUGRA; (B) MSSM without slepton mixing; (C) MSSM with lepton branching ratio set to same as W/Z. For model (B) a limit on the \( \tilde{\chi}_1^\pm \) mass greater than 130 GeV is set (Fig. 1(a)).

DØ has performed four searches \(^2\) with luminosity 1.0-1.1 fb\(^{-1}\) with no excess of data observed. Three mSUGRA-inspired models (with no slepton mixing) are explored: (1) large \( m_0 \) where W/Z decays dominate; (2) 3\( \ell \)-max with slepton mass slightly larger than the \( \tilde{\chi}_2^0 \) mass; (3) heavy squarks where scalar mass unification is relaxed (Fig. 1(b)). For the 3\( \ell \)-max model, a limit of \( M(\chi_1^\pm) > 141 \text{ GeV} \) is found.

2 \( W' \)

Some extensions of the standard model predict the existence of additional, heavy, gauge bosons. DØ has performed a search for a \( W' \) decaying to an electron and neutrino \(^3\) using a dataset of 0.9 fb\(^{-1}\). Data selection requires a high energy electron \( (E_T > 30 \text{ GeV}) \), large missing transverse energy \( (\text{MET} > 30 \text{ GeV}) \) and large transverse mass \( (M_T > 150 \text{ GeV}) \). Figure 2(a) shows the transverse mass distribution (without \( M_T \) cut) for data, background and signal. DØ observes 630 events with an expected background of \( 623 \pm 18^{+83}_{-50} \) events. Therefore, a limit on a \( W' \) mass \( > 965 \text{ GeV} \) is set assuming standard model couplings (Fig. 2(b)).
3 \( Z' \)

CDF has performed a model independent search for narrow resonances decaying to an electron and a positron \(^4\) using 1.3 fb\(^{-1}\) of data. They scan the mass region 150-900 GeV in 4 GeV mass bins looking for an excess of data over predicted background (Fig 3(a)). The small excesses seen are consistent with statistical fluctuations. This is interpreted to exclude a standard model type \( Z' \) with mass below 923 GeV. Additional models are shown in Fig. 3(b).

4 Randall-Sundrum Gravitons

Both CDF and D0 have combined searches in di-electron final states with similar searches in di-photon to explore models of extra dimensions involving Randall-Sundrum gravitons \(^4,5\). Models of extra dimensions attempt to address the hierarchy problem between the strength of the weak force and gravity. At hadron colliders, RS gravitons may be observed in the invariant mass or angular distributions of electron and/or photon pairs. Both experiments observe data in agreement with background predictions and exclude large regions in the graviton mass vs. \( k/M_{pl} \) parameter space (Fig. 4). At \( k/M_{pl}=0.1 \) CDF(D0) exclude gravitons with masses below 889(865) GeV.

5 Excited Electrons

Some models predict that quarks and leptons are composite particles composed of smaller pieces. These models allow for excited quark/lepton states. D0 has carried out a search for excited electrons (\( e^* \)) from the process \( p\bar{p} \rightarrow ee^* \rightarrow ee\gamma \). After selecting events with \( p_T(e_1,e_2,\gamma) > 25,15,15 \) GeV, 259 events are observed with an expected background of 232 \( \pm 3 \pm 29 \) events. From this, limits are set on the mass of the excited electron and the compositeness scale (Fig. 5). For \( \Lambda = 1 \), the limit is \( m_{e^*} > 756 \) GeV. If decays via contact interaction are neglected, D0 finds a limit of \( m_{e^*} > 946 \) GeV for \( \Lambda = m_{e^*} \).
Figure 3: (a) Distribution of di-electron invariant mass with data shown as points with errors and background as the histograms. (b) CDF limit on the cross-section times branching ratio for a spin 1 object along with various models of $Z'$ production.

Figure 4: Limits on extra dimensions using the Randall-Sundrum model from (a) CDF and (b) D0. Limits are set on the parameters $k/M_{Pl}$ and the graviton mass.
6 Neutral, Long-lived Particles

D0 has performed a search for neutral, long-lived particles decaying to two muons after traveling at least 5 cm from the production point. A sample of pair production of neutralinos with R-parity violating decays and long lifetime is used to model the signal. Background is estimated from data to be $0.75 \pm 1.1 \pm 1.1$ events. No events are observed with a decay length in the transverse plane of 5-20 cm. Limits are set on the production cross-section times branching ratio as well as a comparison with a previous result from NuTeV\textsuperscript{7} using the sample model (Fig. 6). This comparison limits the possible interpretations of NuTeV’s result.

7 Summary

The CDF and D0 collaborations have performed numerous searches for new phenomena using leptonic final states. Recent results place limits on associated chargino and neutralino production, extra gauge bosons, Randall-Sundrum gravitons, excited electrons and neutral, long-lived particles. Most of these are the world’s best limits.

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5. D0 Conference Note 5195, \url{http://www-d0.fnal.gov/Run2Physics/WWW/results/np.htm}. 

Figure 5: Limits on excited electrons from D0. (a) shows the limit on the cross-section times branching ratio as a function of the mass of the excited electron. (b) shows the limit on the compositeness scale vs. the mass of the excited electron.
Figure 6: Limits on the cross-section times branching ratio for neutral, long-lived particles decaying to two muons as a function of the lifetime. The area above the (red) line is excluded at the 95% CL. The dark blue shaded region is a 99% CL from D0. The yellow region shows the limit from NuTeV converted to pp collisions at $\sqrt{s} = 1.90$ TeV. The light blue region shows the area favored by a signal interpretation of NuTeV's result.

SEARC\HES IN PHOTON AND JET STATES

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We present recent results from the Collider Detector at Fermilab (CDF) and D\O\ experiments using data from proton-antiproton collisions with $\sqrt{s} = 1.86$ TeV at Run II of the Fermilab Tevatron. New physics may appear in events with high transverse momentum objects, including photons and quark or gluon jets. The results described here are of signature-based searches and model-based searches probing supersymmetry, leptoquarks, $4^{\text{th}}$ generation quarks, and large extra dimensions.

1 Signature Based Searches

1.1 Searches in $\gamma\gamma + \{e, \mu, \gamma, E_T\}$

The signature-based search for two photons plus an additional photon, electron, muon, or missing transverse energy ($E_T$) at CDF in Run II\(^1\) is motivated by an anomalous Run I event\(^2\) containing $ee\gamma E_T$, and the fact that new physics may appear in events with multiple high transverse momentum ($p_T$) particles. This signature may occur in numerous scenarios beyond the standard model (SM): pair-production of supersymmetry (SUSY) particles that yield multiple photons as they decay to lower mass states (for example, through $\chi^0_1 \rightarrow \gamma\chi^0_2$); pair-production and radiative decays of other excited states ($X^* X^* \rightarrow \gamma X\gamma X$); $b'$ quark pair production and decay; fermiphobic Higgs decays yielding photons; and gauge-mediated supersymmetry breaking (GMSB) scenarios. Using data corresponding to an integrated luminosity of $1.0 \text{ to } 1.2 \text{fb}^{-1}$, the number of observed events (4,1,3,0) in each of the four final states ($\gamma\gamma\gamma, \gamma\gamma E_T, \gamma\gamma e, \gamma\gamma\mu$) is consistent with the expected number of events due to SM processes (2.2, 0.24, 6.8, 0.7).
Table 1: Number of expected and observed events for the $\ell\gamma + \{\ell, \gamma, E_T\}$ search, and the dominate SM processes.

<table>
<thead>
<tr>
<th>Expected (SM)</th>
<th>$ee\gamma$</th>
<th>$\mu\mu\gamma$</th>
<th>$e\mu\gamma + X$</th>
<th>$\gamma\gamma$</th>
<th>$\mu\gamma\gamma$</th>
<th>$e\gamma E_T$</th>
<th>$\mu\gamma E_T$</th>
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<tr>
<td>Dominant Source</td>
<td>$Z\gamma$</td>
<td>$Z\gamma$</td>
<td>$Z\gamma$</td>
<td>$e$ fakes $\gamma$</td>
<td>jet fakes $\gamma$</td>
<td>$W\gamma$</td>
<td>$W\gamma$</td>
</tr>
<tr>
<td>Observed</td>
<td>53</td>
<td>21</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>96</td>
<td>67</td>
</tr>
</tbody>
</table>

Figure 1: (left) Cross section limit for GMSB scenario as a function of SUSY breaking scale, assuming prompt photons. (right) Exclusion regions for a GMSB scenario with delayed photons.

1.2 Searches in $\ell\gamma + \{\ell, \gamma, E_T\}$

An additional signature-based analysis at CDF searches for events containing an electron or muon ($\ell = e$ or $\mu$) with $p_T > 25\,\text{GeV}/c$, a photon with $E_T > 25\,\text{GeV}$, and an additional object, which can be an electron, muon, photon, or $E_T^\gamma$ \(^3\). This search, which uses 929 pb\(^{-1}\), is motivated by a 2.7$\sigma$ excess found in Run I in the $\ell\gamma E_T$ final state\(^4\). In that search, 16 events were observed compared to an expectation of $7.6 \pm 0.7$ events due to SM sources. The current search uses the same event selection criteria, but about 10 times the amount of data, and observes 163 $\ell\gamma E_T$ events compared to a SM expectation of $150.6 \pm 13.0$ events. Table 1 shows the expected and observed number of events for the individual channels, as well as the dominant SM process which yields each signature. In all cases, the observations are consistent with the SM.

2 Supersymmetry

2.1 Gauge Mediated Supersymmetry Breaking in $\gamma\gamma E_T$

The spectrum of particles predicted by SUSY theory, which is a popular SM extension, depends upon the mechanism by which SUSY is broken. One possibility is the GMSB scenario, where $\bar{\chi}^0_1\chi^0_1 \rightarrow \gamma \tilde{G} \tilde{G}$, yielding a signature of two photons and $E_T$ from the gravitinos (\tilde{G}). A search using 760 pb\(^{-1}\) at DO expects $2.1 \pm 0.7$ events from SM sources and observes 4 events\(^5\). Assuming short neutralino lifetime, so that the photons are produced promptly, 95% C.L. limits are set at $m(\tilde{\chi}^0_1) > 120\,\text{GeV}/c^2$ and $m(\tilde{\chi}^0_1) > 220\,\text{GeV}/c^2$. Figure 1(left) shows the cross section limit as a function of the SUSY breaking scale parameter $\Lambda$. The other parameters of the model are fixed at the so-called SPS8 values\(^6\).
2.2 Search for Delayed Photons

The detection of photons produced after a delay with respect to the primary interaction may indicate the presence of new phenomena. These photons may stem from decays of heavy, neutral, long-lived particles, such as in the GMSB scenario, with $\chi_1^0 \rightarrow \gamma \tilde{G}$. Using 570 pb$^{-1}$, CDF searches for delayed photons in the $\gamma E_T +$jet channel. The measurement uses a new timing system in the electromagnetic calorimeter. The signal selection requires the photon time of arrival to be $2 < t < 10$ ns, reducing backgrounds from beam halo, cosmic rays, and prompt photons. We expect $1.3 \pm 0.7$ SM events and observe 2 events. Figure 1(right) shows the resulting exclusion regions in the plane of neutralino lifetime and mass, assuming SPS8 parameters as in section 2.1.

2.3 Search for Stopped Gluinos

In split-SUSY, the SUSY scalars are heavy (possibly of the grand unification scale) compared to the SUSY fermions. Split-SUSY models predict long-lived gluinos ($\tilde{g}$), which could have time to hadronize, lose momentum through ionization, and come to rest in the calorimeter. Using 350 pb$^{-1}$, DØ searches for $\tilde{g} \rightarrow \tilde{g} \chi_1^0$ in the jet+$E_T$ channel, where the gluon yields a jet and the unobserved $\chi_1^0$ results in $E_T$. The selection requires one jet with $E > 90$ GeV, in an otherwise empty event. Background from cosmic ray muons is reduced by requiring wide jet showers in the calorimeter with no associated muon, since jets are wide for the hypothetical signal but narrow for muons. Additional beam-induced background is successfully removed. Cross section limits are set as a function of jet energy and translated into gluino mass limits. The search sets a limit of $m(\tilde{g}) > 270$ GeV/$c^2$ for an example with $m(\chi_1^0) = 50$ GeV/$c^2$. Figure 2(left) shows cross section limits as functions of gluino mass.

2.4 Search for Squarks and Gluinos

In the minimal supergravity (mSUGRA) formulation of SUSY, pair produced squarks and gluinos ($\tilde{q}\tilde{q}, \tilde{g}\tilde{g}, \tilde{q}\tilde{g})$ can decay through $\tilde{q} \rightarrow q\chi_1^0$ and $\tilde{g} \rightarrow q\tilde{q}\chi_1^0$ to yield signatures with 2, 3, or 4 jets from the quarks and $E_T$ from the neutralinos. A search at DØ combines these three channels, categorized by the number of jets, to set 95% C.L. limits on the squark and gluino...
masses of $m(\bar{q}) > 375\text{ GeV}/c^2$ and $m(\tilde{g}) > 290\text{ GeV}/c^2$, using 960 pb$^{-1}$. Figure 2(right) shows the excluded regions in the squark–gluino mass plane.

3 Leptoquarks

3.1 3rd Generation Vector Leptoquarks

Leptoquarks may provide a link between the families of quarks and leptons at higher mass scales. They appear in a variety of beyond the SM theories. At the Fermilab Tevatron, leptoquarks could be predominately pair produced through quark anti-quark annihilation. Using 322 pb$^{-1}$ of CDF data, we consider the case of 3rd generation vector leptoquarks ($VLQ_3$), where each decays exclusively to $\tau b$. The decay products, $\tau b\bar{b}$, are reconstructed as two jets from the $b$ quarks, one leptonically decaying tau, and one hadronically decaying tau. At 95% C.L., we set a limit of $m(VLQ_3) > 317\text{ GeV}/c^2$ for a model with Yang-Mills type couplings. Figure 3(left) shows the cross section limits as a function of $VLQ_3$ mass for two models.

3.2 3rd Generation Scalar Leptoquarks

Complementary to the search mentioned above, leptoquarks may also appear as spin=0 objects. A search at DØ assumes 3rd generation scalar leptoquarks ($SLQ_3$) are pair produced and decay primarily via $SLQ_3 \rightarrow \nu b$, but also allows for $SLQ_3 \rightarrow \tau t$. The considered decay products, $\nu b\tau\bar{b}$, yield a signature of 2 $b$-jets and $E_T$ from the neutrinos. Based on 310 pb$^{-1}$, the 95% C.L. limit is $m(SLQ_3) > 213\text{ GeV}/c^2$. If $SLQ_3 \rightarrow \tau t$ is explicitly forbidden, the limit is $m(SLQ_3) > 219\text{ GeV}/c^2$. Figure 3(right) shows these results.

3.3 2nd Generation Scalar Leptoquarks

As with the other leptoquark searches, to avoid producing flavor changing neutral currents, it is common to assume that generation number is conserved in the decays. Therefore a search for 2nd generation scalar leptoquarks ($SLQ_2$) at DØ assumes the decays contain muons or muon neutrinos, and second generation quarks ($q = s$ or $c$). We consider the branching ratio $\beta \equiv B(SLQ_2 \rightarrow \mu q) = 0.5$, so that pair production and decay leads to a final state of $\mu q\mu q$. The signature is one muon, 2 jets, and $E_T$. In 1 fb$^{-1}$, 6.1 $\pm$ 1.1 events are expected from SM
Figure 4: (left) Cross section limit for $2^{\text{nd}}$ generation scalar leptoquarks as a function of mass. (right) Observed limit and theoretical prediction for a $4^{\text{th}}$ generation $b'$ model, in the search for new particles that couple with $Z$ bosons.

backgrounds, while 6 events are observed. The 95% C.L. limit is $m(SLQ2) > 214 \text{GeV}/c^2$, as shown in figure 4(left).

4 Search for $b'$

A search for $b'$ 4\textsuperscript{th} generation quarks is carried out at CDF, as a specific case to quantify the acceptance, in a general search for new particles that couple to $Z$ bosons and jets\textsuperscript{15}. The production and decay chain considered in the analysis is $qq \rightarrow g \rightarrow b'b' \rightarrow b'ZbZ$, with one $Z$ decaying through $Z \rightarrow \ell\ell$ and the other decaying through $Z \rightarrow q\bar{q}$. The leptonically decaying $Z$ is reconstructed in the $Z \rightarrow e\mu$ and $Z \rightarrow \mu\mu$ channels, with lepton $p_T > 20 \text{GeV}/c$, and we require $\geq 3$ jets with $E_T > 30 \text{GeV}$. The leading background, from SM $Z$+jet events, is estimated by using events with $< 3$ jets to predict the contribution for events with $\geq 3$ jets. Using $1.1 \text{fb}^{-1}$, the 95% C.L. limit is $m(b') > 270 \text{GeV}/c^2$, as indicated by the intersection of the observed and theoretical curves in figure 4.

5 Large Extra Dimensions

Large extra dimensions have been proposed\textsuperscript{16} as a potential solution to the hierarchy problem between the weak and gravitational forces. In these models, only gravitons ($G$) can propagate in the $n$ extra dimensions of the $4+n$ dimensional bulk of space-time. The Planck scale ($M_{\text{Planck}}$) is so large, and the strength of gravity so small, due to the large extent ($R$) of the extra spacial dimensions. To address the hierarchy problem, the new effective Planck scale ($M_D$) is related to the other quantities through $M_{\text{Planck}}^{2} \sim R^{n} M_{D}^{2+n}$. The gravitons can be directly produced via $qq \rightarrow gG$, $gg \rightarrow qG$, or $gg \rightarrow gG$, in all cases yielding an energetic jet from the quark or gluon, and $E_T$ from the graviton. The measurement at CDF, using $1.1 \text{fb}^{-1}$, observes 779 events while expecting $819 \pm 17$ events from SM sources. This is an update, using the same selection but more data, of the previously published results\textsuperscript{17}. These new results are used to place lower limits on $M_D$ and upper limits on $R$. The lower limits on $M_D$ are given in figure 5, with corresponding upper limits on $R$ of $R < 0.27 \text{mm}$, $R < 3.1 \times 10^{-6} \text{mm}$, $R < 9.9 \times 10^{-9} \text{mm}$, $R < 3.2 \times 10^{-10} \text{mm}$, and $R < 3.1 \times 10^{-11} \text{mm}$ for $n = 2, 3, 4, 5,$ and 6.
6 Conclusions

Various searches for new physics at DØ and CDF using photons and jets are constraining specific models, while signature based searches are looking for deviations from the SM. Discoveries may be just around the corner, as additional analyses and larger data sets are considered.

Acknowledgments

We thank our colleagues with the CDF and DØ collaborations, the Moriond conference organizers and participants, and the funding agencies for making this work possible. We also thank the European Union Marie Curie Program for an accommodation grant.

References

SUPERSYMMETRY BREAKING MADE EASY, Viable, AND GENERIC

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The kind of supersymmetry that can be discovered at the LHC must be very much flavor-blind, which used to require very special intelligently designed models of supersymmetry breaking. This led to the pessimism for some in the community that it is not likely for the LHC to discover supersymmetry. I point out that this is not so, because a garden-variety supersymmetric theories actually can do this job.

1 Introduction

LHC is coming! It is finally taking us to the energy scale of the weak interaction, $G_F^{-1/2} \approx 300$ GeV, known as an important energy scale for more than seven decades since Fermi's 1933 paper on the nuclear beta decay. It is a historic moment in science and I am very excited to be part of this new era. Whenever physicists had crossed a threshold of studying a new force, it resulted in a big paradigm change. The atomic scale (scale of quantum electrodynamics) led to the revolutionary discovery of quantum mechanics. The nuclear scale (scale of quantum chromodynamics) revealed a new layer of matter and showed the non-perturbative quantum field theory to be essential in our description of nature. We are all looking forward to whatever paradigm change the weak scale will bring us.

However, there has been a growing concern in the community, especially among the theorists, that we may not find anything surprising at the electroweak scale. I have been bitten by this bug, too. The reasoning is very simple. If there is new rich physics below the TeV scale such as supersymmetry and/or extra dimensions, why haven't we seen its impact already on precision electroweak and flavor-physics experiments? Because we haven't seen such impacts, it is unlikely that there is rich new physics below the TeV scale and most likely we will not find anything spectacular at the LHC.

*This talk is based on the collaboration with Yasunori Nomura*
Even though I had plunged into this pessimism myself, I have now completely turned around back to optimism. I do think it is quite likely for the LHC to find something exciting such as supersymmetry. I would like to tell you why I made this 180 degrees change in my attitude.

The issue is the following. Supersymmetry, if present at the TeV scale, must be a broken symmetry because we have not seen any superpartners yet. The problem is that it has been believed that breaking supersymmetry is very difficult, and certainly is not generic among supersymmetric theories. Even among those that do break supersymmetry, they are rather difficult to use for constructing phenomenologically viable models, hence most of them are dead on arrival. A very small fraction of the minority then survive, after an elaborate model-building gymnastic, namely the “alive” theories are “pockets of insurgency” in the barren land. An elaborate model is like a beautiful artwork, intelligently designed, which is what makes its creator(s) proud, but is by definition special and fragile. I can’t stop feeling that Nature is unlikely to rely on such a fragile elaborately built artwork for the foundation of its inner working.

I now feel the situation is very different from the previous perception. There is a large fraction of supersymmetric models that can successfully break supersymmetry in a phenomenologically viable fashion. And it appears rather robust, namely a change in parameters, such as choice of the gauge groups, number of flavors, does not spoil its success. This observation makes it much more plausible that supersymmetry has a broad and robust foundation to be realistic, and makes me feel that it could well be there waiting for us at the LHC.

2 Hierarchy Problem

It is often said that the hierarchy problem has been overemphasized as a reason to expect rich physics at the LHC. There is some truth in it, because, after all, it is an aesthetic problem. However, I’d like to first remind you that the hierarchy problem was actually solved once in the history of physics, and we can draw some lesson from it:3

In the classical electromagnetism, the only dynamical degrees of freedom are electrons, electric fields, and magnetic fields. When an electron is present in the vacuum, there is a Coulomb electric field around it, which has the energy of

$$\Delta E_{\text{Coulomb}} = \frac{1}{4\pi\varepsilon_0} \frac{e^2}{e}. \tag{1}$$

Here, $r_e$ is the “size” of the electron introduced to cutoff the divergent Coulomb self-energy. Since this Coulomb self-energy is there for every electron, it has to be considered to be a part of the electron rest energy. Therefore, the mass of the electron receives an additional contribution due to the Coulomb self-energy:

$$(m_e c^2)_{\text{obs}} = (m_e c^2)_{\text{bare}} + \Delta E_{\text{Coulomb}}. \tag{2}$$
Figure 2: (Left) The Coulomb self-energy of the electron. (Middle) The bubble diagram which shows the fluctuation of the vacuum. (Right) Another contribution to the electron self-energy due to the fluctuation of the vacuum.

Experimentally, we know (now) that the "size" of the electron is small, $r_e \lesssim 10^{-17}$ cm. This implies that the self-energy $\Delta E$ is at least a few GeV, and hence the "bare" electron mass must be negative to obtain the observed mass of the electron, with a fine cancellation like:

$$0.000511 = (-3.141082 + 3.141593) \text{ GeV}. \quad (3)$$

Even setting a conceptual problem with a negative mass electron aside, such a fine cancellation between the "bare" mass of the electron and the Coulomb self-energy appears troublesome. In order for such a cancellation to be absent, Landau and Lifshitz concluded that the classical electromagnetism cannot be applied to distance scales shorter than $e^2/(4\pi\varepsilon_0 m_e c^2) = 2.8 \times 10^{-13}$ cm. This is a long distance in the present-day particle physics’ standard.

The resolution to this problem came from the discovery of the anti-particle of the electron, the positron, or in other words by doubling the degrees of freedom in the theory. The Coulomb self-energy discussed above can be depicted by a diagram Fig. 2, left where the electron emits the Coulomb field (a virtual photon) which is felt (absorbed) later by the electron itself. But now that we know that the positron exists, and we also know that the world is quantum mechanical, one should think about the fluctuation of the "vacuum" where a pair of an electron and a positron appears out of nothing together with a photon, within the time allowed by the energy-time uncertainty principle $\Delta t \sim \hbar/\Delta E \sim \hbar/(2m_e c^2)$ (Fig. 2, middle). This is a new phenomenon which didn’t exist in the classical electrodynamics, and modifies physics below the distance scale $d \sim c\Delta t \sim \hbar c/(2m_e c^2) = 200 \times 10^{-13}$ cm. Therefore, the classical electrodynamics indeed does hit its limit of applicability at this distance scale, much earlier than $2.8 \times 10^{-13}$ cm as was exhibited by the problem of the fine cancellation above. Given this vacuum fluctuation process, one should also consider a process where the electron sitting in the vacuum by chance annihilates with the positron and the photon in the vacuum fluctuation, and the electron which used to be a part of the fluctuation remains instead as a real electron (Fig. 2, right). V. Weisskopf calculated this contribution to the electron self-energy, and found that it is negative and cancels the leading piece in the Coulomb self-energy exactly:

$$\Delta E_{\text{pair}} = \frac{1}{4\pi \varepsilon_0 r_e^2}. \quad (4)$$

After the linearly divergent piece $1/r_e$ is canceled, the leading contribution in the $r_e \to 0$ limit is given by

$$\Delta E = \Delta E_{\text{Coulomb}} + \Delta E_{\text{pair}} = \frac{3\alpha}{4\pi m_e c^2} \log \frac{\hbar}{m_e c r_e}. \quad (5)$$

There are two important things to be said about this formula. First, the correction $\Delta E$ is proportional to the electron mass and hence the total mass is proportional to the "bare" mass of the electron,

$$(m_e c^2)_{\text{obs}} = (m_e c^2)_{\text{bare}} \left[ 1 + \frac{3\alpha}{4\pi} \log \frac{\hbar}{m_e c r_e} \right]. \quad (6)$$

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*Do you recognize $\pi$?*
Therefore, we are talking about the "percentage" of the correction, rather than a huge additive constant. Second, the correction depends only logarithmically on the "size" of the electron. As a result, the correction is only a 9% increase in the mass even for an electron as small as the Planck distance $r_e = 1/M_{Pl} = 1.6 \times 10^{-33}$ cm.

The fact that the correction is proportional to the "bare" mass is a consequence of a new symmetry present in the theory with the antiparticle (the positron): the chiral symmetry. In the limit of the exact chiral symmetry, the electron is massless and the symmetry protects the electron from acquiring a mass from self-energy corrections. The finite mass of the electron breaks the chiral symmetry explicitly, and because the self-energy correction should vanish in the chiral symmetric limit (zero mass electron), the correction is proportional to the electron mass. Therefore, the doubling of the degrees of freedom and the cancellation of the power divergences lead to a sensible theory of electromagnetism applicable to very short distance scales.

In the Standard Model, the Higgs potential is given by

$$V = m^2 |H|^2 + \frac{\lambda}{4} |H|^4,$$

where $\langle H \rangle^2 = -m^2/2 \lambda = (176 \text{ GeV})^2$. Because perturbative unitarity requires that $\lambda \lesssim 1$, $-m^2$ is of the order of $(100 \text{ GeV})^2$. However, the mass squared parameter $m^2$ of the Higgs doublet receives a quadratically divergent contribution from its self-energy corrections. For instance, the process where the Higgs doublets splits into a pair of top quarks and come back to the Higgs boson gives the self-energy correction

$$\Delta m_{top}^2 = -\frac{m^2}{4\pi^2} \frac{1}{r_H},$$

where $r_H$ is the "size" of the Higgs boson, and $h^2 \approx 1$ is the top quark Yukawa coupling. Based on the same argument in the previous section, this makes the Standard Model not applicable below the distance scale of $10^{-17}$ cm, according to the Landau–Lifshitz criterion. This is the hierarchy problem.

The motivation for supersymmetry is to make the Standard Model applicable to much shorter distances so that we can hope that the answers to many of the puzzles in the Standard Model can be given by physics at shorter distance scales. In order to do so, supersymmetry repeats what history did with the positron: doubling the degrees of freedom with an explicitly broken new symmetry. Then the top quark would have a superpartner, the stop, whose loop diagram gives another contribution to the Higgs boson self energy

$$\Delta m_{stop}^2 = \frac{m^2}{4\pi^2} \frac{1}{r_H^2}.$$  

The leading pieces in $1/r_H$ cancel between the top and stop contributions, and one obtains the correction to be

$$\Delta m_{top}^2 + \Delta m_{stop}^2 = -\frac{m^2}{4\pi^2} \left( \frac{m^2}{m_t^2} - m_t^2 \right) \log \frac{1}{r_H^2 m_t^2}. $$

One important difference from the positron case, however, is that the mass of the stop, $m_t$, is unknown. In order for the $\Delta m^2$ to be of the same order of magnitude as the tree-level value $m^2 = -2\lambda v^2$, we need $m_t^2$ to be not too far above the weak scale. To $V$ stop mass is already a fine tuning at the level of a percent. Similar arguments apply to masses of other superpartners that couple directly to the Higgs doublet. This is the so-called naturalness constraint on the superparticle masses.

It is worth pondering if the mother nature may fine-tune. Now that the cosmological constant appears to be fine-tuned at the level of $10^{-120}$, should we be worried really about the fine-tuning of $v^2/M_{Pl}^2 \approx 10^{-30}$? In fact, some people argued that the hierarchy exists because intelligent
life cannot exist otherwise. On the other hand, a different way of varying the hierarchy does seem to support stellar burning and life. We don’t get into this debate here, but we’d like to just point out that a different fine-tuning problem in cosmology, horizon and flatness problems, pointed to the theory of inflation, which in turn appears to be empirically supported by data. We just hope that proper solutions will be found to both of these fine-tuning problems and we will see their manifestations at the relevant energy scale, namely TeV.

3 Why We Were Pessimistic

Supersymmetry or not, we expect some interesting physics to appear below TeV scale if the hierarchy problem is to be avoided by some stabilization mechanism. The problem is that it is difficult to understand why we have not seen its impact on flavor-changing neutral currents especially in the beautiful $B$ physics data and electroweak precision measurements. Are we on the wrong track to naively hope that the stabilization mechanism is just around the corner? Or is there rather a good reason why it doesn’t show its fingerprints despite our best detective work? This is a question that applies to any candidate physics beyond the standard model at the TeV scale.

The problem with supersymmetry is well-known, having been discussed already for a several decades. In the Minimal Supersymmetric Standard Model (MSSM), which is the supersymmetric extension of the Standard Model with the smallest particle content, there are staggering 107 additional parameters beyond the nineteen in the Standard Model. And if you throw dice in this huge parameter space, you almost always end up with a parameter set that is already excluded by the data. For example, the off-diagonal elements in the mass-squared matrices must be less than a few per mill of the mass eigenvalues for three types of squark and two types of slepton mass matrices. Also the mixing between the scalar partners of left- and right-handed fermions need to be very much identical to the mixing among the fermions. Overall, the probability of “hitting” the phenomenologically viable parameter sets would be down by a product of many factors of hundreds. Why are the unwanted parameters small, supersymmetry-breaking effects flavor-blind? Unless there is a good reason, the whole idea of sub-TeV supersymmetry to stabilize the hierarchy appears a remote chance.

In addition, breaking supersymmetry appears difficult and highly non-generic among supersymmetric theories. Known models require a specific choice of gauge groups and matter content, often together with a special choice of the superpotential terms with a global symmetry imposed; global symmetries are usually regarded unlikely in a fundamental theory of quantum gravity such as string theory.

There are several popular mechanisms to achieve flavor-blind supersymmetry breaking: gauge mediation,10 gaugino mediation,11 and anomaly mediation.12 The supersymmetry-breaking effects are “mediated” to the supersymmetric standard model via gravity or gauge interactions, guaranteeing their flavor-blindness. Even though these mechanisms do work, my problem has been that the models must be written in a very careful and elaborate fashion. Small changes in the models, such as a different choice of the gauge group or matter content, tend to destroy their success, such as restoring supersymmetry, allowing for flavor-dependent effects, destabilizing the vacuum. My feeling has been that they do not represent a likely choice by Nature.

Let us see how careful and elaborate models are needed with an example. Gauge mediation, at the first sight, is a beautiful idea. It has a set of messenger particles $f$ and $\tilde{f}$ that carry standard-model quantum numbers. Once supersymmetry is broken by a vacuum expectation value of field $X$, the messengers do not have a supersymmetric spectrum; the masses of the bosons are split from those of the fermions. Then their loops induce masses for squarks, sleptons, and gauginos (Fig. 3, left). So far so good.

However, one has to ask the question how the supersymmetry-breaking expectation value is
generated for $X$. It requires a separate gauge theory with a rather complicated particle content and specific potential, on which a global symmetry is imposed to make sure that it breaks supersymmetry. Only a small fraction of supersymmetric models serve this purpose. This gauge theory is coupled to the messengers also in a constrained specific way through yet another gauge interaction and singlet particles (Fig. 3, right). Overall, we need nearly decoupled three “boxes” to make the whole mechanism possible.

If a realistic model of nature has to rely on a carefully constructed elaborate mechanism, even though it is logically possible, I am not sure if that is Mother Nature’s choice. In some sense, we rely on supersymmetry to avoid fine-tuning in electroweak symmetry breaking and flavor, but we end up fine-picking or intelligently designing a model. Even though it is a philosophical point and not very scientific, I would be very much happier if we can do without fine-picking. If we think Nature doesn’t fine-tune, she probably doesn’t fine-pick either; she is way smarter than us, after all.

4 The New Generic Scheme

My main point here is that I actually do not need to be very intelligent to achieve flavor-blind supersymmetry breaking. Pretty much the dumbest supersymmetric extension of the standard model would do it.

We still need the messengers, non-chiral particles coupled to the standard model. Such particles are known to arise generically in string theory, and people have been trying hard to get rid of such junk. In addition, again generically, one expects other gauge groups, with their own quarks; more junk. Most of them tend to be non-chiral. I claim these junk are precisely what we need. No fine-picking.

We do not impose a global symmetry on the model. We write the most general potential consistent with the gauge symmetries. The lowest dimension operator that couple the messengers and the other quarks is precisely what brings the supersymmetry-breaking effects to the messengers, and hence to the standard model in a flavor-blind way. The gauge groups can be pretty much anything, $SU(N_c)$, $SO(N_c)$, $Sp(N_c)$; all classical groups. (Let us focus on $SU(N_c)$ case below but the other groups do the same thing.) The parameters (e.g., the gauge couplings, mass scales, etc) do not need to be tuned against each other. They have to satisfy certain inequalities, such that the possibly flavor-sensitive Planck-scale physics would be subdominant compared to the gauge-mediated contributions, or the number of flavors in the other “QCD” is in the range $N_c < N_f < \frac{3}{2}N_c$. No fine-tuning here.

---

4For instance, the first three-generation models from heterotic strings based on Tian-Yau manifold has six vector-like families, and an extra $E_8$ for disposal. The only significant requirement here is that there should better not be chiral particles that couple to both the standard model and the other gauge groups. I thank Mirjam Cvetic on this point.
How this simple scheme works requires a little technical discussion. For this range of the $N_f$, the other "QCD" becomes strong at some energy scale $\Lambda$, and the low-energy limit is known to be described by yet another gauge theory $SU(N_f - N_c)$, whose quarks $q$ and $\bar{q}$ couple to the mesons $S_{ij} = \bar{Q}_i Q_j$ ($i,j = 1, \cdots, N_f$) of the original "QCD" through the potential

$$ W = m^{ij}_Q S_{ij} - \bar{S}_{ij} q^i q^j. \tag{11} $$

We assume $m^{ij}_Q \ll \Lambda$. Then this potential does not have a solution to the supersymmetric minimum

$$ \frac{\partial W}{\partial S_{ij}} = m^{ij}_Q - \bar{q}^i q^j = 0 \tag{12} $$

because $m^{ij}_Q$ has rank $N_f$ while $\bar{q}^i q^j$ rank $N_f - N_c$. This supersymmetry-breaking minimum is actually a local minimum, but the tunneling to the global supersymmetric minimum has an exponentially long lifetime. The lowest dimension operator for the messengers become

$$ M \bar{f} \bar{f} + \frac{\Lambda}{M_{Pl}} \langle S \rangle \bar{f} \bar{f}, \tag{13} $$

exactly what is needed for the gauge mediation (Fig. 3, left).

In fact, any models that break supersymmetry can be used the same way. Just the vector-like "junka" coupled to the standard model, and the lowest dimension operator to link them.

This, I believe, is a good news for string theory. String theory is now believed to have many many solutions, some $10^{500}$ of them. The vast majority of them have huge cosmological constants and do not resemble our universe; they do not support life and do not get observed by scientists like us. We have lost very many solutions by this cut. Getting standard model is another severe cut on the numbers. It would be nice if a large fraction of the remaining solutions would lead to successful supersymmetry breaking and phenomenologically viable models; then Nature may well have given us one. The simple scheme presented here suggests that a significant fraction of the remaining solutions indeed may well do so.

Experimental consequences are pretty much the same as the phenomenology of gauge-mediated models people have been discussing in the literature. The dark matter particle is the gravitino. Even though it is a bad news for direct detection experiments, it opens up an interesting possibility of producing gravitinos of spin $3/2$ at colliders. There may be extra photons or long-lived charged particles in the supersymmetry events. The mass spectrum of superparticles tell us about the quantum numbers of the messengers even though they are beyond the reach of direct production. Finally, the linear-collider precision of superparticles may reveal the presence of the light particles in the other "QCD".
5 Conclusions

Even though supersymmetry is a beautiful idea to solve the fine-tuning or hierarchy problem, we will not see it at the LHC unless it comes out flavor-blind. Theorists used to be a kind of control freak to write special models that ensure the flavor-blindness of supersymmetry. This fine-picking of models made us uncomfortable, feeling the chance for its discovery at the LHC remote. But after some more thoughts, it turned out that we don’t really need fine-picking to break supersymmetry in a flavor-blind fashion. It is easy to write a model, which is phenomenologically viable, and the scheme is very generic, a kind of spectrum one expects from the string theory. Pretty much the dumbest extension of the supersymmetric standard model would do the job.

I do not share anymore the spreading concern that LHC is not likely to discover exciting physics because we have not seen any hints of it yet. Quite generically in the “landscape” of supersymmetric theories, we expect the superparticles to come out flavor-blind and therefore well hidden from the current beautiful data. I suspect this is probably not specific to supersymmetry. More thoughts may well reveal why we have not seen hints of TeV-scale new physics yet, even though it is waiting to be discovered at the LHC.

Acknowledgments

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ELECTROWEAK STUDIES AND SEARCHES IN INCLUSIVE HIGH $Q^2$ $eP$
COLLISIONS

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Recent results from the analysis of inclusive high $Q^2$ collisions of longitudinally polarised electrons or positrons and protons recorded at HERA are presented. The data are used to measure electroweak parameters and to search for new physics beyond the standard model.

1 Introduction

The HERA machine and the two collider experiments H1 1 and ZEUS 2 have been taking data since 1992. Since 2003 both experiments have upgraded interaction regions including strong focusing magnets, and the instantaneous luminosity has increased by a factor 3. Furthermore, spin rotators have been added near H1 and ZEUS, leading to collisions of protons with longitudinally polarised electrons or positrons. The machine operates at a centre-of-mass energy $\sqrt{s} = 320\text{ GeV}$. The inclusive cross-sections presented here are measured in two reactions: neutral current (NC) and charged current (CC) interactions. For NC interactions, the scattered electron is detected in the tracking systems and the calorimeter. In contrast, CC interactions involve a scattered neutrino, which escapes detection. The kinematics of the scattering process is determined by the negative square of the momentum transfer, $Q^2$, the proton momentum fraction of the struck quark Bjorken, $x$, and the inelasticity, $y$. These observables are related by $Q^2 = sxy$, so only two of them are independent.

1.1 The Neutral Current Cross-Section

The NC inclusive cross-sections may be written in terms of the polarisation dependent generalised structure functions $\tilde{F}_2$, $x\tilde{F}_3$, $\tilde{F}_L$,

$$\frac{d^2\sigma^{NC}}{dx dQ^2} = \frac{2\pi\alpha^2}{xQ^4} \left[ Y_+ \tilde{F}_2(x, Q^2) + Y_- x\tilde{F}_3(x, Q^2) - y^2 \tilde{F}_L(x, Q^2) \right].$$

The $\mp$ sign in front of the $x\tilde{F}_3$ term corresponds to the case of $e^+p$ scattering. The helicity functions are defined as $Y_{\pm} = 1 \pm (1 - y)^2$. The contribution from $\tilde{F}_L$ is small and can be neglected at high $Q^2$. The generalised structure functions $\tilde{F}_2$ and $x\tilde{F}_3$ can be expressed as:

$$\tilde{F}_2 = F_2 - (v_e \pm P_e a_e)\kappa \frac{Q^2}{Q^2 + M_Z^2} F_2^{\gamma Z} + (v_e^2 + a_e^2 \pm P_e 2a_e a_e)\kappa^2 \left[ \frac{Q^2}{Q^2 + M_Z^2} \right]^2 F_2^{\gamma Z},$$

$$x\tilde{F}_3 = -(a_e \pm P_e v_e)\kappa \frac{Q^2}{Q^2 + M_Z^2} xF_3^{\gamma Z} + 2a_e v_e \pm P_e (v_e^2 + a_e^2)\kappa^2 \left[ \frac{Q^2}{Q^2 + M_Z^2} \right]^2 xF_3^{\gamma Z}.$$
Figure 1: Single-differential cross-sections for $e^+p$ scattering in neutral current and charged current reactions.

The structure function $F_2$ corresponds to the case of pure photon exchange. The electron couplings to the $Z^0$ are denoted by $a_e$ and $v_e$. $P_e$ is the lepton polarization and $\kappa^{-1} = \sin^2(2\theta_W)$ is related to the weak mixing angle. The structure functions $F_2^{Z^0}$ and $x^{F_3^{Z^0}}$ correspond to $\gamma/Z^0$ interference effects. The terms related to pure $Z^0$ exchange are small in the kinematic domain accessible at HERA. In the Born approximation the structure functions can be decomposed into parton density functions (PDFs)

$$
\left[ F_2, F_2^{Z^0}, F_2^{\gamma Z} \right] = \sum_q x \left[ e_q^2, 2e_q v_q, (v_q^2 + a_q^2) \right] (q + \bar{q}),
$$

$$
\left[ F_3^{Z^0}, F_2^{\gamma Z} \right] = \sum_q x [2e_q a_q, 2v_q a_q] (q - \bar{q}).
$$

Here $e_q$, $v_q$, and $a_q$ are the quark electric charge and couplings to the $Z^0$ boson. The quark and anti-quark densities of the proton are given by $q$ and $\bar{q}$. The sums run over the five quark flavours $q = \{u, d, s, c, b\}$ accessible at HERA.

1.2 The Charged Current Cross-Section

At Born level the CC cross-section can also be expressed in terms of the parton density functions as follows:

$$
\frac{d^2\sigma^{CC}}{dx dQ^2} (e^+p \to e^+X) = (1 + P_e) \frac{1}{x} \frac{G_F^2 M_W^4}{4\pi (Q^2 + M_W^2)^2} \left[ (1 - y)^2 (xd + xs + xb) + (x\bar{d} + x\bar{s} + x\bar{b}) \right],
$$

$$
\frac{d^2\sigma^{CC}}{dx dQ^2} (e^-p \to e^-X) = (1 - P_e) \frac{1}{x} \frac{G_F^2 M_W^4}{4\pi (Q^2 + M_W^2)^2} \left[ (xu + xc) + (1 - y)^2 (x\bar{d} + x\bar{s} + x\bar{b}) \right].
$$

Here $G_F$ is the Fermi constant and $M_W$ is the mass of the W boson.
Figure 2: Double-differential cross-sections for $e^\pm p$ scattering in neutral current reactions and the resulting measurement of $xF_3$, presented as a function of $x$ for fixed $Q^2$ values as denoted in the figure.

2 Single-Differential Unpolarised Cross-Sections

The single-differential cross-sections of the NC and CC reactions as a function of $Q^2$ have been measured by H1 \cite{3,4} and ZEUS \cite{5,6,7}. The results are shown in Figure 1. Compared to HERA I, the precision of the $e^-p$ data has improved significantly. At low $Q^2$, the NC cross-section is dominated by the photon propagator and behaves like $1/Q^2$. In contrast, the CC cross-section is almost constant at low $Q^2$. At high $Q^2$ the NC and CC cross-section are of similar size, a direct way to visualise the electroweak unification at scales near the heavy gauge boson masses. Significant differences between the $e^+p$ and $e^-p$ NC cross-sections are due to the $\gamma/Z^0$ interference. For the CC process, the differences between $e^+p$ and $e^-p$ scattering is due to the different types of valence quarks entering the reaction and from the different helicity factors for the valence quarks.

3 Double-Differential Unpolarised Cross-Sections

The double-differential NC cross-section has been measured for $e^-p$ and $e^+p$ collisions. Data from HERA II are combined with HERA I data. The ZEUS data are shown in Figure 2. The difference between the $e^-p$ and $e^+p$ datasets is used to extract \cite{8} the structure function $xF_3$. In order to enhance the statistical precision, the data points are all transformed to $Q^2 = 1500 \text{GeV}^2$. The combined H1 and ZEUS results on the structure function $xF_3^{\gamma Z}$ are also shown in Figure 2. Using the $xF_3$ data, the sum rule $\int_{0.02}^{0.65} \frac{xF_3}{x} dx = \frac{5}{4}$ can be tested at high $Q^2$. The experimental result

$$\int_{0.02}^{0.65} \frac{xF_3}{x} dx = 1.21 \pm 0.09(\text{stat}) \pm 0.08(\text{sys})$$

is in agreement with expectations from QCD fits and with the sum rule, when extrapolated to $x = 0$ and $x = 1$. 
Figure 3: The cross-section integrated up to $y = 0.9$ for $e^+p$ scattering in charged current reactions as a function of the lepton polarisation, $P_e$.

Table 1: Right-handed charged current cross-sections derived from the polarisation dependence of the total charged-current cross-section.

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<thead>
<tr>
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<th>$\sigma_{CC}$ (extrapolated to $P_e = 1$)</th>
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<tbody>
<tr>
<td>$e^-p$ H1 (prel.)</td>
<td>$-0.9 \pm 2.9_{\text{stat}} \pm 1.9_{\text{sys}} \pm 2.9_{\text{pol}}$</td>
</tr>
<tr>
<td>$e^+p$ ZEUS (prel.)</td>
<td>$0.8 \pm 3.1_{\text{stat}} \pm 5.0_{\text{sys+pol}}$</td>
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<th>$\sigma_{CC}$ (extrapolated to $P_e = -1$)</th>
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<tbody>
<tr>
<td>$e^-p$ H1</td>
<td>$-3.9 \pm 2.3_{\text{stat}} \pm 0.7_{\text{sys}} \pm 0.8_{\text{pol}}$</td>
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</tbody>
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4 Polarised Cross-Sections

4.1 The polarised charged current cross-section

The charged current cross-section predicted by the standard model has a simple linear dependence on the polarisation $P_e$ of the incoming lepton

$$\sigma^{CC}(P_e) = (1 \pm P_e) \sigma^{CC}(P_e = 0).$$

Beyond the standard model, there may be a non-vanishing contribution from right-handed charged currents, for example by the exchange of a heavy $W_R$ boson. Figure 3 shows the integrated charged current cross-section for $Q^2 > 400$ GeV and $y < 0.9$ as a function of the lepton polarisation, $P_e$. The data are in good agreement with a linear dependence of $\sigma^{CC}$ on $P_e$. If the cross-section is extrapolated to $P_e = +1$ ($P_e = -1$) for incoming electrons (positrons), limits on right-handed charged currents can be set. The limits are summarised in Table 1. If translated to a limit on a massive $W_R$ boson with right-handed couplings of electroweak strength, then $W_R$ masses smaller than 208 GeV are excluded at 95% confidence level (using only the H1 result on polarised $e^+p$ data).
for $e^\pm p$ collisions. Here $P_R > 0$ and $P_L < 0$ denote the lepton polarisation $P_e$ reached by HERA for different setting of the spin-rotators. The combined results are shown in Figure 4. The asymmetry $A^\pm$ reaches values of 0.4 for high values of $Q^2$. The measurements are in good agreement with standard model predictions based on fits of unpolarised data and they provide direct evidence for parity violation mediated by electroweak effects in NC reactions at high $Q^2$. The asymmetry has a different sign depending on the lepton charge. $A^+$ and $A^-$ are of similar size, as expected from the $Z^0/\gamma$ interference term. Note that the asymmetry is sensitive to the quark vector couplings.
5 Electroweak Fits

QCD parameters and parton densities have been determined with high precision from HERA data\textsuperscript{9,10}. Recently, new fits have been performed where in addition to QCD parameters some electroweak parameters are determined\textsuperscript{11,12}. Examples of fits to the $W$ mass and the $u$ quark axial and vector couplings are shown in Figure 5. The precision on the $W$ mass is not competitive to direct measurements at LEP and Tevatron. Still, it is interesting to see that the CC propagator mass measured in $e^+e^-$ collisions at high $Q^2$ is consistent with on-shell measurements performed at other machines. The determination of the $u$ and $d$ quark couplings to the $Z^0$ improve significantly over the measurements performed at other colliders.

6 Searches for New Physics

6.1 Leptoquarks

Leptoquarks (LQs) are bosons with both leptonic and baryonic quantum numbers. A new quantum number $F = 3B + L$ is introduced. In a generic model developed by Buchmuller Rückl and Wyler\textsuperscript{13}, LQs are grouped by their electroweak quantum numbers and 14 different types of LQs are accessible at HERA. LQs can be pair-produced in $e^+e^-$ or $p\bar{p}$ collisions. At HERA, LQs can be singly produced from the incoming lepton and a quark from the proton. If the LQ mass is smaller than the HERA centre-of-mass energy, then the LQ is produced on-shell and its mass can be reconstructed using the formula $M = \sqrt{sxy}$. For LQ masses beyond the HERA centre-of-mass energy, the $Q^2$ spectrum is sensitive to the presence of LQs. Inclusive data from deep inelastic scattering at HERA I and HERA II have been searched for LQs\textsuperscript{14,15,16} and no signal was found. The reconstructed LQ mass and mass-dependent limits on LQ production determined from HERA II data are shown in Figure 6 for one particular type of LQ. The limits on LQ production depend on the LQ type and mass. For masses of 200 GeV the limit on the LQ coupling is of order 0.02. For couplings of electromagnetic strength ($\lambda = 0.3$) the mass limit is of order 300 GeV.
6.2 Contact interactions

The HERA I and HERA II data have been searched for contact interactions, large extra dimensions and finite quark radii $^{17,18}$. These effects can modify the single-differential cross-sections at high $Q^2$. No evidence of such phenomena was found and limits have been set. Figure 7 shows the single-differential NC cross-sections as a function of $Q^2$ normalised to the SM prediction. The expectation from a specific contact interaction model is also shown. Limits are set at 95% confidence on a broad variety of models, as shown in Figure 7. Depending on the model, compositeness scales $\Lambda$ up to 7.5 TeV are excluded. For models with large extra dimensions, the HERA I-II limit on the mass scale is $M_S > 0.88$ TeV. The limit on the quark radius for point-like electrons is found to be $R_q < 0.67 \times 10^{-16}$ cm.
7 Summary

The HERA collider is a unique facility to measure inclusive $e^\pm p$ reactions at high momentum transfer with polarised leptons. The electroweak structure function $x F_3$ is extracted from neutral current $e^\pm p$ reactions. The polarisation dependence of charged current reactions provides an interesting method to look for right-handed charged currents which do not exist in the standard model. For neutral current reactions, the polarisation asymmetry gives direct evidence for parity violating effects at high momentum transfer $Q^2$. Fits to the inclusive cross-section measurements are made in order to simultaneously extract QCD and electroweak parameters. The HERA results on the $u$ and $d$ quark axial and vector couplings complement the LEP results on heavy quarks. Finally, the inclusive cross-sections are searched for physics beyond the standard model in leptoquark production and contact interactions. No evidence for new physics is found.

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EXCLUSIVE SEARCHES AT HERA

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The latest results from the H1 and ZEUS collaborations on exclusive searches for rare Standard Model processes and physics Beyond the Standard Model are presented. Intriguing events containing high transverse momentum leptons are observed by H1. The ZEUS searches for similar event topologies are also presented. Competitive limits on excited neutrinos and lepton flavour violation are set by the H1 collaboration. A generalised search by H1 for events containing high \( P_T \) objects is also presented.

1 Introduction

Both the H1 and ZEUS experiments pursue a comprehensive physics programme at the HERA \( ep \) collider. In addition to measurements of the structure of the proton, deep inelastic scattering at HERA at a centre-of-mass energy of up to 318 GeV provides an ideal environment for the study of rare processes, for setting constraints on the Standard Model (SM) and for searches for new particles and physics Beyond the Standard Model (BSM).

2 Rare Standard Model Processes

2.1 Events Containing Isolated Leptons and Missing Transverse Momentum

Events containing a high transverse momentum (\( P_T \)) isolated electron or muon and large missing transverse momentum have been observed at HERA \(^1\). The dominant SM mechanism for producing such topologies is the production of real \( W \) bosons, \( e p \rightarrow e W^{\pm} X \), with subsequent leptonic decay \( W \rightarrow l\nu \). An excess in the HERA I (1994-2000) data sample was reported by the H1 collaboration \(^2\), whereas such an excess was not confirmed by the ZEUS collaboration in a similar search \(^3\).

The H1 analysis has been updated to include new \( e^\pm p \) data from the ongoing HERA II (2003-2007) running period \(^4\), resulting in a total analysed luminosity of 442 pb\(^{-1}\). A total of 56 events are observed in the data, compared to a SM prediction of 54.9 ± 7.6. The hadronic transverse momentum (\( P_T^X \)) spectra of the \( e^\pm \) data are presented in figure 1. At large values of \( P_T^X (> 25 \text{ GeV}) \) 18 \( e^+ p \) data events are observed compared to the SM prediction of 7.8±1.3. The observed excess is equivalent to a fluctuation of approximately 3\( \sigma \). The excess in not observed in the \( e^- p \) analysis.

A similar analysis has been performed by the ZEUS collaboration on 432 pb\(^{-1}\) of \( e^\pm \) taken in the 1996-2000 and 2003-2007 running periods \(^5\),\(^6\). The SM expectations of the ZEUS search are similar to the H1 search, however the ZEUS analysis does not observe an excess in \( e^+ p \)
Figure 1: The hadronic transverse momentum spectra of the observed events in the H1 isolated lepton analysis. The sample is divided into $e^-p$ (shown on the left for $\mathcal{L} = 184$ pb$^{-1}$) and $e^+p$ (right for $\mathcal{L} = 258$ pb$^{-1}$) data samples. The data are shown as points, the filled histogram is the SM expectation and the shaded band is the total SM error. The signal component, dominated by real $W$ production, is shown by the hatched histogram.

Figure 2: The efficiency of the H1 and ZEUS isolated lepton searches for $ep \rightarrow eWX$ (left $W \rightarrow e\nu_e$, right $W \rightarrow \mu\nu_\mu$) events with $P_T^X > 25$ GeV as a function of the lepton polar angle $\Theta_{e,\mu}$. The arrows in the figures denote the H1 and ZEUS candidate events. The generated spectrum for $ep \rightarrow eWX$ Monte Carlo is also shown.

data; with only one electron candidate observed in the $P_T^X > 25$ GeV region compared to a SM prediction of $1.5 \pm 0.2$ and three muon candidates observed compared to a SM prediction of $3.1 \pm 0.5$. The ZEUS data also agrees well with the SM predictions for the $e^-p$ data with 5(2) electrons (muons) observed compared to the expectation of $3.8 \pm 0.6$ ($2.2 \pm 0.3$) events.

H1 and ZEUS have formed a combined working group to clarify differences between the searches of the two collaborations. Comparisons of the efficiencies of the two searches as a function of $P_T^X$ show that the ZEUS and H1 have similar efficiency for selecting $e^\pm p \rightarrow e^\pm WX$ events at high values of $P_T^X$. The more restricted phase space of the ZEUS search leads to a slightly lower efficiency than the H1 search. It can be seen in figure 2 that in regions of the lepton polar angle $\Theta_{e,\mu}$ common to both searches the efficiencies of the two searches are similar. H1 and ZEUS candidate events are illustrated by arrows on the figures, it can be seen that the majority of the H1 candidates lie within the ZEUS search acceptance, indicating that the ZEUS search is also sensitive to physics which could be responsible for the H1 excess.

2.2 Multi-lepton Events

Searches for multi-electron and multi-muon production at high transverse momentum have been previously carried out by the H1 experiment\(^8\). The most recent H1 search was performed using 275 pb$^{-1}$ data taken during the HERA I and HERA II running periods. The main SM process for multi-lepton production in $ep$ collisions is photon-photon interaction $\gamma\gamma \rightarrow l^+l^-$. 


where quasi-real photons radiated from the incoming electron and proton interact to produce a pair of leptons. The H1 search studied $ee, \mu\mu, ee, eee$ and $e\mu\mu$ topologies, searching for events containing at least two high $P_T$ leptons. The data are found to be overall in good agreement with the SM. Events are observed at masses greater than 100 GeV in $e^+p$ data.

The ZEUS collaboration has recently made a new search for multi-electron events using 296 pb$^{-1}$ $e^+p$ data from the HERA I and HERA II running periods. The invariant mass ($M_{12}$) distributions for the two highest $P_T$ electrons in the $ee$ and $eee$ topologies are shown in Figure 3, where overall good agreement with the SM expectations is observed. $M_{12} > 100$ GeV events are observed in both the $ee$ and $eee$ topologies; the number of such events observed is in agreement with the SM expectation.

3 Searches for Physics Beyond the Standard Model

3.1 General Search for New Phenomena

A model-independent search for deviations from the SM was previously performed by H1, using HERA I data. This search has been repeated by H1 using 159 pb$^{-1}$ of $e^+p$ data. High $P_T$ final state configurations involving electrons ($e$), muons ($\mu$), jets ($j$), photons ($\gamma$) or neutrinos ($\nu$) are considered. All final state configurations containing at least two such objects with $P_T > 20$ GeV in the central region of the detector are investigated and classified into exclusive event classes, $\mu - j, e - j\nu, j - j - j$ etc.

The yields for the different classes are shown in Figure 4; data events are found in 21 event classes and good agreement is observed between data and the SM expectations in most event classes. A non-biased statistical method is used to search for deviations of the data with respect to the SM in the summed scalar transverse momentum ($\sum P_T$) and total invariant mass ($M_{\text{all}}$) distributions sensitive to new physics. Good agreement is observed in all event classes. The distribution of the probability of obtaining a result at least as different from the SM as the obtained spectrum ($P$) for scans of $\sum P_T$ in the different event classes is also shown in Figure 4. The data distribution is consistent with the SM.

3.2 Excited Neutrinos

The fermion mass hierarchy can naturally be explained if the SM fermions are composite, in which case excited states may exist. A minimal extension of the SM can incorporate
excited fermions such as excited neutrinos ($\nu^*$). Considering only electroweak interactions, the excitation part of the lagrangian is

$$L_{\nu^*}\nu = \frac{1}{2\Lambda} T_R T^R \sigma^{\mu\nu} \left[ g f \frac{\tau^\nu}{2} \delta_{\mu\nu} W^\nu_{\nu} + g' f' \frac{Y}{2} \delta_{\mu\nu} B^\nu_{\nu} \right] F_L + \text{h.c.},$$

where the new weights $f$ and $f'$ multiply the SM coupling constants $g$ and $g'$ corresponding to the weak SU(2) and electromagnetic U(1) sectors respectively. The corresponding gauge boson fields are denoted by $W$ and $B$. The matrix $\sigma^{\mu\nu} = (i/2)[\gamma^\mu, \gamma^\nu]$, $\tau$ are the Pauli matrices, and $Y$ is the hypercharge. The compositeness scale $\Lambda$, together with $f$ and $f'$, determines the production cross section. Excited neutrinos can be produced in $ep$ collisions at HERA via the $t$-channel charged current (CC) reaction $e^+ p \rightarrow \nu^* X$. The cross section is much larger in $e^- p$ collisions than in $e^+ p$ collisions due to helicity enhancement specific to CC-like processes. H1 have recently presented a new search on 114 pb$^{-1}$ of $e^- p$ data, a dataset almost an order of magnitude larger than previously published HERA analyses.

Excited neutrinos are searched for in the decay channels $\nu^* \rightarrow \nu \gamma$, $\nu^* \rightarrow eW$ and $\nu^* \rightarrow \nu Z$. The $W$ and $Z$ bosons are reconstructed in the hadronic decay channel. No signal was observed in any channel. Limits on the production cross section are calculated and translated into exclusion limits in the plane ($f/\Lambda, M_{\nu^*}$), under the assumption $f = f'$ or $f = -f'$ (Figure 5). For $f = -f'$ (maximal $\gamma\nu\nu^*$ coupling) and assuming $f/\Lambda = 1/M_{\nu^*}$, $\nu^*$s with masses less than 188 GeV are excluded at 95% CL. The results confirm the unique sensitivity at HERA to excited neutrinos with masses beyond the reach of LEP.
3.3 Leptoquark Production and Lepton Flavour Violation

Particle interactions in the SM conserve lepton flavour, although there is no underlying symmetry supporting this feature. Experimental evidence for lepton flavour violation (LFV) in solar and atmospheric neutrino oscillations has been reported\textsuperscript{19,20}. The experimental bounds on neutrino masses imply very small LFV effects in the charged lepton sector. The observation of such effects would clearly indicate new phenomena beyond the SM.

At HERA, LFV processes $e \mu \rightarrow \mu X$ or $e \tau \rightarrow \tau X$ lead to final states with a muon or tau and hadronic system $X$. The LFV process can proceed via the exchange of a leptoquark (LQ), a boson with both lepton and baryon quantum number. Leptoquarks arise naturally as a colour triplet scalar or vector boson in many extensions of the SM such as grand unified theories, supersymmetry, compositeness and technicolor. The H1 collaboration has recently performed a new search for LFV phenomena using 66.5 pb$^{-1}$ $e^{\pm}p$ data and 13.7 pb$^{-1}$ $e^{-}p$ data\textsuperscript{21}.

The LFV processes $e \mu \rightarrow \mu X$ and $e \tau \rightarrow e \tau X$ can be attributed to LQs produced predominantly by electron-quark fusion. LQs are classified into 14 types with respect to the quantum numbers spin, isospin and chirality within the framework of the Buchm{"u}ller-R{"u}ckl-Wyler (BRW) model\textsuperscript{22}. Leptoquarks carry both lepton (L) and Baryon (B) quantum numbers. The fermion number $F = L + 3B$ is assumed to be conserved, taking values of $F = 2$ for $e^{-}q$ processes and $F = 0$ for $e^{+}q$ processes. For LQ masses ($m_{LQ}$) well below the $e^{\pm}p$ centre-of-mass energy, the $s$ channel production of $F = 2$ ($F = 0$) LQs in $e^{-}p$ ($e^{+}p$) collisions dominates. For $m_{LQ}$ greater than the centre-of-mass energy the $s$ and $u$ channel processes become of equal importance. The BRW model assumes lepton flavour conservation such that the LQs produced in $ep$ collisions decay only to $eX$ or $\nu_{e}X$ final states. A general extension of the BRW model allows for the decay of LQs to final states containing a lepton of different flavour ($\mu$ or $\tau$) and a jet. Non-zero couplings to an electron-quark pair and to a muon(tau)-quark pair are assumed.

In the H1 analysis, high $P_{T}$ muon and tau signatures are searched for. In the muon case, the signature is an isolated high $P_{T}$ muon back-to-back to the hadronic system in the transverse plane. In the tau case electron, muonic and hadronic tau decays are all considered. In the case of electronic tau decays, large missing momentum is expected in the muonic decay the signature is very similar to the muon case and in the hadronic decay the signal topology is a di-jet event with no leptons. The tau-jet is characterised by a narrow energy deposit in the calorimeter and low track multiplicity. The tau-jet candidates are selected via a neural net algorithm\textsuperscript{23}.

Only one candidate event was found, in the $e^{+}p$ data, in the hadronic $\tau$ decay channel. This observation is in agreement with the SM and no evidence for LFV was found. Limits on couplings to 14 different LQs as a function of $m_{LQ}$ are derived. An example of the limits under the assumption that the tau- and electron-first generation quark couplings are equal is shown in Figure 6. The H1 results are directly comparable to previous limits set by ZEUS\textsuperscript{24} and are found to be similar. At hadron colliders LQs are mainly produced in pairs independently of coupling, thus searches cannot constrain LFV couplings. Lower mass limits on the second (third) generation leptoquarks extend up to 250 GeV (150 GeV), depending on the assumptions made\textsuperscript{25,26}. Similarly lower mass bounds from $e^{+}e^{-}$ annihilation reach values of 100 GeV\textsuperscript{27}.

4 Summary

Many searches for new physics have been performed at HERA by the H1 and ZEUS collaborations. No evidence for lepton flavour violation or $\nu$'s is observed. Interesting events containing isolated leptons and missing $P_{T}$ as well as the multiple high $P_{T}$ leptons at high masses are observed by H1. H1-ZEUS working groups have been formed to clarify the high $P_{T}$ lepton excess observed by H1 and to extend the reach of BSM searches at HERA to its uttermost.
Figure 6: Limits on the coupling constants $\lambda_{\chi_1} = \lambda_{\chi_2}$ as a function of the leptoquark mass $m_{\chi}$ for $F=0$ scalar (left) and vector (right) leptoquarks. Regions above the lines are excluded at 95% CL. The notation $q_1$ indicates that only processes involving first generation quarks are considered.

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THE STANDARD MODEL HIGGS BOSON AT LHC: RECENT DEVELOPMENTS

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The search for the Higgs boson is one of the most important tasks of two multi-purpose experiments: CMS and ATLAS. Both experiments are expected to start taking data in 2008, when the LHC begins its operation at 14 TeV. The most recent results from the CMS and ATLAS Standard Model Higgs boson discovery potential are presented, obtained using the most recent full simulation and reconstruction software as well as with detailed treatment of systematic uncertainties. Results show that the whole Standard Model Higgs boson mass range will be covered with about \( \sim 10^{-5} \) fb\(^{-1}\).

There are four main experiments at the Large Hadron proton-proton Collider (LHC) at CERN\(^2\), two of them are general purpose detectors: CMS (Compact Muon Solenoid)\(^3\) and ATLAS\(^4\) (A Toroidal Lhc ApparatuS). One of the most important tasks for these experiments is to search for the Higgs boson, the last missing piece of the Standard Model (SM) that is still not confirmed experimentally.

1 The Standard Model Higgs boson production and decay

The main production channels for the SM Higgs boson are: gluon fusion via top loop, vector boson (W/Z) fusion, W/Z associated production, \(t\bar{t}\) fusion. The production cross section at the LHC as a function of the the Higgs boson mass is shown in Fig. 1\(^5\) (right). In the Higgs boson search analyses, NLO corrections for signal and background processes were used when available. At LHC, 14 TeV proton-proton collider, the dominant production mechanism is gluon fusion, whereas at Tevatron, proton-anti-proton 2 TeV collider, the most relevant mechanism is W/Z associated production, see e.g. Ref. \(^6\).
The Higgs decay modes (see Fig. 1, left plot) can be divided into two different mass ranges. For $M_H \leq 135\text{GeV}/c^2$ the Higgs boson mainly decays into $b\bar{b}$ and $\tau^+\tau^-$ pairs with branching ratios of about 85% and 8% respectively. Both $b\bar{b}$ and $\tau^+\tau^-$ Higgs boson decay modes (as well as modes into $c\bar{c}$ and gluon pairs) are [almost] impossible for discovery due to overwhelming background level. The most useful Higgs decays in this mass range at the LHC is the decay into photon pairs. It reaches a branching fraction of up to $2 \times 10^{-3}$. Its importance lies in the clean environment and controllable background, which makes it an important discovery channel for small Higgs boson masses.

For Higgs masses above $135 \text{GeV}/c^2$ the main decay modes are those into $WW$ and $ZZ$ pairs. These are very important discovery channels, in particular $H \rightarrow ZZ^{(*)} \rightarrow 4\mu$ is called the “golden mode” for its very clean (4 muons) final state, which provides a low/medium integrated luminosity level for discovery of the Higgs boson for all possible Higgs boson masses in consideration ($\sim 115 - 600\text{GeV}/c^2$).

2 Searches for the SM Higgs boson

This section discusses the detailed analyses available for many of the SM Higgs boson production and decay channels from both CMS and ATLAS collaborations. They include:

- inclusive $H \rightarrow \gamma\gamma$\textsuperscript{7,8},
- inclusive $H \rightarrow WW/ZZ$\textsuperscript{9,10},
- $qqH$ production mode with Higgs decaying into $\tau\tau, \gamma\gamma, WW/ZZ$\textsuperscript{11,12},
- $W/Z+H$ production mode with Higgs decaying into $\gamma\gamma, WW$\textsuperscript{13},
- $t\bar{t}H$ production mode with Higgs decaying into $b\bar{b}, \gamma\gamma$\textsuperscript{14,15}.

Figure 1: Left: Higgs production cross sections at the LHC for the various production mechanisms as a function of the Higgs mass. The full QCD-corrected results for the gluon fusion $gg \rightarrow H$, vector-boson fusion $qq \rightarrow VV$, $qq \rightarrow Hq$, vector-boson bremsstrahlung $q\bar{q} \rightarrow V^* \rightarrow HV$ and associated production $gg, q\bar{q} \rightarrow Ht\bar{t}$ are shown. Right: Branching ratios of the dominant decay modes of the SM Higgs particle. All relevant higher-order corrections are taken into account.
2.1 Inclusive $H \rightarrow \gamma\gamma$

The distinct feature of this analysis is the assignment of a weight to every event depending on the "quality" of reconstructed photon candidates. One can find the signal-background spectrum for different photon categories in Fig. 2\textsuperscript{16}. There is a clear gain in exploiting this feature: getting the most for significance from events with less background contamination. This analysis also took advantage of optimizing the signal vs. background significance via Neural Network techniques, which decreased the integrated luminosity level needed for 5$\sigma$ discovery roughly by factor of 2.

2.2 Inclusive $H \rightarrow ZZ$

In the $H \rightarrow ZZ^{(*)} \rightarrow 4\mu$ analysis counting and Log-Likelihood ratio techniques were used as well as genetic algorithm techniques\textsuperscript{17} to optimize the outcome of the analysis, which led to $M_{4\mu}$-dynamic cuts and about 1$\sigma$ improvement in terms of signal to background significance all over the $M_{4\mu}$ spectrum. Effort was made to make analysis as close as possible to "as one would do it with real data": lepton reconstruction and isolation techniques from data were developed and tested.

A detailed study of systematic uncertainties was performed in this and many other analyses by CMS and ATLAS. Some of the main systematic uncertainties come from the major background cross sections (for many analyses) and NLO k-factors as well as possible contribution from unaccounted backgrounds (including processes, which become important when fake leptons, for example, are taken into account). Also estimations were made to re-scale the local excess of significance due to the fact that we look for a narrow resonance in a broad range of $M_H$.

2.3 Inclusive $H \rightarrow WW$

In $H \rightarrow \gamma\gamma$ analysis, one uses sidebands of huge, but smooth background in $M_{\gamma\gamma}$-spectrum; the, inclusive $H \rightarrow WW$ analysis uses control regions in phase space where signal "contamination" is minimized (see Fig. 3\textsuperscript{16}, top). Signal and irreducible WW background events were also re-weighted with dynamic NLO k-factors (for $p_T^{WW}$ distribution) in order to properly calculate the discovery potential for this important channel (the procedure is proposed in Ref.\textsuperscript{18} and it improves significance considerably).

2.4 Vector boson fusion $H \rightarrow WW$

As in the previous one, this analysis is also a counting experiment and uses signal and control parts of phase space for the discovery region and background control (see Fig. 3\textsuperscript{12}, bottom). The advantage in this analysis is the ability to tag two additional forward jets.

2.5 Discovery reach summary

Figure 4 (left)\textsuperscript{16} shows the signal significance as a function of the Higgs boson mass for 30 fb$^{-1}$ of integrated luminosity for the different Higgs boson production and decay channels with CMS. Similar results for ATLAS are shown on Fig. 4 (right)\textsuperscript{12}.

Figure 5 (left)\textsuperscript{16} shows the integrated luminosity needed for the 5$\sigma$ discovery of the inclusive Higgs boson production $pp \rightarrow H + X$ with the Higgs boson decay modes $H \rightarrow \gamma\gamma$, $H \rightarrow ZZ \rightarrow 4\ell$, and $H \rightarrow WW \rightarrow 2\ell2\nu$ - the three front runners among all possible decay channels. H $\rightarrow \gamma\gamma$ dominates the discovery reach up to $m_H \sim 130$GeV/$c^2$, similarly $H \rightarrow WW \rightarrow 2\ell2\nu$ dominates for $\sim 150 - 180$ GeV and $H \rightarrow ZZ \rightarrow 4\ell$ - in the rest of the possible Higgs boson masses. Figure 5 (right) shows the Higgs boson mass measurements precision, $\Delta M_H/M_H$, as a function of the
Higgs boson mass for 30 fb$^{-1}$ of integrated luminosity for the different Higgs boson production and decay channels with CMS.

3 Summary

Both collaborations, CMS and ATLAS, show that discoveries may be expected already at integrated luminosity, about 1 fb$^{-1}$. The SM Higgs boson, if it exists, is expected to be discovered by the time we reach $L \approx 10$ fb$^{-1}$. By the time we get $L \approx 30$ fb$^{-1}$ we should be able to measure the Higgs boson mass with precision of about 0.1%.

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Figure 3: Inclusive $H \rightarrow WW$ analysis. Top: The angle between the leptons in the transverse plane is shown for the signal and the different backgrounds and a luminosity of 10 fb$^{-1}$, left plot is for the signal cuts and right plot is for the WW background normalization region where all signal cuts are applied, except the one on the lepton invariant mass. VBF $H \rightarrow WW$ analysis. Bottom: Distribution of the transverse mass $M_T$ for Higgs boson masses of 160 GeV/c$^2$ is shown (left - is for "signal-like" part of the phase space, and right - is for "background-like" part of the phase space).

Figure 4: Left: The signal significance as a function of the Higgs boson mass for 30 fb$^{-1}$ of the integrated luminosity is shown for the different Higgs boson production and decay channels for analyses with CMS. Right: The corresponding summary plot for ATLAS is shown. The difference in reach curves mainly comes from NLO vs. LO cross sections and full simulation vs. fast simulation for CMS vs. ATLAS results. (There is an update expected for ATLAS official results.)
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CHALLENGES FOR EARLY DISCOVERY IN ATLAS AND CMS

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The challenges for a discovery of new physics with 1 fb\(^{-1}\) of LHC data for ATLAS and CMS are discussed. Four specific examples are chosen: a deviation of QCD jet distributions at high \(E_T\), high-mass dilepton pairs, Higgs search in the WW decay channel, and low mass supersymmetry.

1 Introduction

The Large Hadron Collider (LHC) is a proton-proton collider with a center-of-mass energy of 14 TeV, currently under construction at CERN, Geneva. At the time of this conference\(^a\), the LHC was still planning to have an engineering run in the fall of 2007, in order to establish single beam operation at the injection energy of 450 GeV, and provide first collisions at fairly low luminosity, at 900 GeV center-of-mass energy\(^b\). The first collisions at \(\sqrt{s} = 14\) TeV could occur in spring 2008, and with a steadily increasing luminosity during the next 26 weeks of proton-proton running, the ATLAS and CMS experiments might collect an integrated luminosity close to 1 fb\(^{-1}\) by the end of 2008. In order to be able to present first results at Moriond 2009, ATLAS and CMS face a number of challenges, some of which will be discussed in this paper. Hereby we will focus on four examples of possible signs of new physics in first data.

The first challenges that ATLAS and CMS face are still severe: the completion of subdetector construction, installation and commissioning, establishing reliable detector operation, and getting the trigger, data acquisition and detector calibration infrastructure to work as designed. The data is to be distributed and analyzed on the Grid, and this also needs to be commissioned.

ATLAS is currently undergoing such a (Monte Carlo) data generation and analysis challenge, which will also lead to performance estimates that update the physics TDR (technical design report) from 1999\(^1\). CMS has produced an updated physics TDR in the summer of 2006\(^2\).

At any energy, soft (low \(p_T\)) hadronic interactions, or “minimum bias” events, will be most common. Various Monte Carlo generators differ significantly in their predictions of the cross

\(^a\)Talk given at Moriond 2007, Electroweak Interactions and Unified Theories, March 10-17, 2007
\(^b\)However, delays in the schedule now make this engineering run increasingly unlikely.
section and the charged particle multiplicity at $\sqrt{s} = 14$ TeV. Since these minimum bias events set the background for trigger and reconstruction, measuring their properties will be first priority for ATLAS and CMS. Furthermore, these events are useful for tracking studies and calorimeter intercalibrations.

2 QCD (di)jets at high $E_T$

The large center-of-mass energy of the LHC gives access to a kinematic region of QCD jet production that has never been probed before. With 1 fb$^{-1}$ of data, jets of 3-3.5 TeV transverse energy ($E_T$), and di-jet masses of up to 6 TeV are accessible. New physics, in the form of quark substructure and excited quarks, contact interactions, or resonances, could appear, as shown in Figure 2.

Two main uncertainties must be tackled before any excess can be interpreted as new physics: parton density function (pdf) uncertainties, and the jet energy scale.

![Figure 2: Left: deviations from the QCD prediction of the dijet mass distribution, for two hypothetical models of new physics, each at three different masses. The vertical bars represent the statistical error with 1 fb$^{-1}$ of data. Center: exclusion and discovery limits on the cross section for various new particles, as a function of the particle mass. Right: significance for an exclusion or a discovery of contact interactions, for 1 fb$^{-1}$ of data, as a function of the (inverse of the) contact interactions scale. Plots from CMS.](image)

The pdf uncertainties arise mainly from the uncertainties on the gluon distribution function at high $x$. The high $E_T$ jet data from ATLAS and CMS can in fact be used to constrain these uncertainties. In order not to sweep new physics under the rug in such a procedure, it is important to fit also complementary processes, like photon plus jet production, in pdf fits, and to measure jet production over a large kinematic range: new physics is often more central, whereas pdf effects show up over all phase space. ATLAS has shown that with 1 fb$^{-1}$ of data, already an improvement from the current situation can be obtained, as shown in Figure 3. Beyond that, however, the systematic errors on the data must be tackled: reducing these errors help more than adding more luminosity. These systematic uncertainties concern mainly the jet energy scale.

The jet energy scale will initially be known only from test beam, cosmics and calibration systems, to not better than 5-10%. With first data, a data-driven jet energy scale determination program must be started immediately. This is a major effort. As a comparison: it has taken the DØ collaboration five years of continuous effort by a large group of people to reach now 2% uncertainty on the energy scale of jets between 30 and 200 GeV. ATLAS and CMS aim for 2-3% after one year, using photon plus jets, Z plus jets, and top-quark pair events. The latter uses the W mass constraint on the two light-quark jets, the former two need a calibration of the electromagnetic energy scale first. Jets from b-quarks need a different calibration than light-quark jets.
Concerning new physics, CMS estimates a discovery potential with 1 fb$^{-1}$ of data for excited quarks up to masses of 3.4 TeV, for diquarks in an $E_6$ model up to 3.7 TeV, and for the scale of contact interactions up to 7.7 TeV. The sensitivity of ATLAS is expected to be similar.

3 High mass lepton pairs

New physics may well show up in high-mass lepton pairs. Resonances such as $Z'$, gravitons in the Randall-Sundrum model, or Kaluza-Klein excitations in models with universal extra dimensions may be seen in the dilepton mass distribution, on top of a Drell-Yan continuum background that falls rapidly with increasing mass. Large extra dimensions of the ADD-model type will not show up as resonances, but as deviations from the mass and angular spectrum of Drell-Yan lepton pairs.

Reconstruction of electrons and muons requires a well-calibrated and aligned detector. The alignment of the CMS and ATLAS inner detectors will be a significant task. CMS has some 20000 silicon modules, or a total of 120000 alignment parameters to be determined; ATLAS has about 6000 modules. Initially, alignment will come from survey measurements, dedicated hardware systems, and cosmic-ray muon data. The ATLAS silicon strip detector includes a system of frequency scanning interferometers, that measure the movements of whole structures (barrels, discs) very precisely over timescales of hours; the CMS inner detector has a laser calibration system. The ATLAS muon spectrometer has an extensive set of laser alignment systems. Eventually, tracks from the data will provide alignment information, but it will be a challenge to do this for all modules, with little residual systematics.

The momentum scale of muons is very much determined by the alignment accuracy of the inner detector and the muon system. It is calibrated with muon pairs from decay of resonances such as the Z boson, and $J/\psi$ and $Y$. Three days of data taking at $10^{30}$ cm$^{-2}$ s$^{-1}$ will provide a sample of more than $10^5$ $Z \rightarrow \mu^+ \mu^-$ events, and eventually the muon momentum scale can be determined to better than 0.1%.

Electrons are measured both in the tracking system and in the electromagnetic calorimeter (ECAL). The CMS ECAL consists of lead tungstate crystals, and will in first instance be intercalibrated to a 0.4 to 2.0% uniformity with single electrons and minimum bias events. The ATLAS ECAL is a lead and liquid argon calorimeter; the goal is a 0.4 to 1.0% uniformity. Both detectors will derive the overall energy scale from $Z \rightarrow e^+ e^-$ events to better than 0.1%.

It should be noted that the energy scale is affected by many effects (magnetic field uniformity, material in the detector, response, alignment etc), and that all these need to be disentangled. After all, the energy scale needs to be extrapolated from the Z peak to high $p_T$, where the new
Figure 4: Muon-pair mass distribution from the decay of a 1 TeV Z’, with ideal alignment (left), or alignment as expected at LHC start-up (center) (also denoted short term alignment) in CMS. Right: track pT resolution for a pT = 100 GeV track in CMS, as a function of pseudorapidity η, for ideal, short-term and long-term (more than a few fb−1 of data) alignment.

physics is expected. A careful Monte Carlo modeling is crucial.

Trigger and reconstruction efficiencies must be obtained from the data itself, and it is important to include redundant and as-little-bias-as-possible triggers in the trigger menu. Also redundant object reconstruction methods are needed: e.g. muons must be reconstructed in the inner detector, in the calorimeter, and in the muon system, so that efficiencies and fake rates can be determined.

On the theoretical side, there are still uncertainties on the Standard Model prediction, originating from missing higher orders in the calculation, scale variations, and pdf uncertainties. ATLAS and CMS will try to select control samples from data to measure the Standard Model background, but one should realize that new physics will probably show up in “untypical” corners of phase space. Therefore, good Monte Carlo predictions are still important, and some NLO calculations are still needed.

As shown in Figure 5, a Z’ could be discovered with 1 fb−1 up to masses of 2-2.8 TeV, depending on the Z’ couplings, Randall-Sundrum graviton resonances up to 2.3 TeV, and ATLAS and CMS would be sensitive to a 6 TeV fundamental Planck scale if there would exist three large extra dimensions of the type of the ADD model (4 TeV for six extra dimensions).

Figure 5: Left: integrated luminosity needed for a Z’ discovery in CMS, as a function of Z’ mass, in various models with different Z’ couplings. Center: significance of a discovery of large extra dimensions in the ADD model for various luminosities, for three extra dimensions, as a function of the fundamental Planck scale M_p. Right: CMS reach in the mSUGRA parameters m0 and m1/2, for tan β = 10, A = 0, μ > 0, for 1 fb−1 of data.

4 Higgs search in the WW → ℓℓνν channel

One of the most promising Higgs search channels with early data is the search for a Higgs in the decay channel H → WW → ℓℓνν, for a Higgs boson in the mass range 150-170 GeV. CMS
estimates that a discovery can be made with less than 1 fb\(^{-1}\) of data\(^4\), and ATLAS also considers this, in particular for Higgs production through WW fusion, the most promising channel.

Since there are two neutrinos in the final state, the mass resolution is poor, and this is essentially a counting experiment. Therefore, it is extremely important to have a good understanding of the background. The major backgrounds, WW and top-quark pair production, are extracted from the data itself through a procedure involving several control samples. The uncertainties related to this procedure have been studied in detail, and seem to be under control. A small, but important, component to WW production comes from gluon-gluon fusion processes.

5 Supersymmetry

Supersymmetry (SUSY) is a theoretically attractive candidate for physics beyond the Standard Model, and there are several arguments in favor of supersymmetry at the TeV scale, accessible at the LHC. If the LHC does not find any evidence for supersymmetry, it is extremely unlikely that supersymmetry is the answer to open issues in the Standard Model like the hierarchy problem. Also the interpretation of the lightest supersymmetric particle as a dark matter candidate will be problematic, certainly if direct detection and astroparticle physics experiments also do not find evidence.

If the SUSY mass scale is only just above the Tevatron limits, SUSY is the prime candidate for an early discovery at the LHC. A search for SUSY in early data must be robust (able to cope with background uncertainties and a non-optimal detector) and general, yet efficient. Excellent opportunities exist in final states with high \(E_T\) jets, and significant missing transverse energy (\(\not{E}_T\)). In order to suppress the QCD background and facilitate triggering, it is possible to further demand one (or more) high \(p_T\) lepton(s). Such final states typically arise in the decay chains of squarks and gluinos, which will be copiously produced at the LHC, if kinematically allowed.

Assuming R-parity conservation, the lightest SUSY particle will escape the detector unseen; if R-parity is violated there will be little \(\not{E}_T\), but still events with many high \(E_T\) jets. Standard Model backgrounds to the SUSY search mainly come from top quark pair production, Z or W boson production in association with jets, and QCD jet production. An example of the \(\not{E}_T\) distribution in signal and background is shown in Figure 6 (left).

![Figure 6: Left: \(\not{E}_T\) distribution in multi-jets final states, for potential SUSY signals at 0.5, 1.0 and 1.5 TeV mass scales, and \(\not{E}_T\) in various backgrounds. Center: \(\not{E}_T\) distribution of Z (\(\rightarrow \nu\nu\)) plus at least two jets, and how this background to SUSY could be estimated from Z (\(\rightarrow \mu\mu\)) events, properly scaled. Right: Expected signal for top-quark pairs in 100 pb\(^{-1}\) of early ATLAS data, without using b-tagging. Shown is the invariant mass distribution of the three jets assigned to a hadronically decaying top quark.](image)

There are several challenges for an early SUSY discovery: reconstruction of leptons, jets, and \(\not{E}_T\) in busy events; fake \(\not{E}_T\) from detector effects; trying to be general, yet efficient, in the SUSY search; and understanding the Standard Model background well.

The \(\not{E}_T\) is a measure of energy carried away by escaping, unmeasured particles, such as neutrinos, or (semi)stable weakly interacting particles in new physics models. Measuring \(\not{E}_T\)
relies on accurate reconstruction of the transverse momentum balance, and is affected by detector effects (holes, noise, punchthrough), and a finite resolution. QCD jets can have real $E_T$ due to neutrinos, and it is a challenge to understand the high $E_T$ tail well.

Reconstruction of objects in a busy environment can be tested at the LHC in top-quark pair production events, which constitute an excellent calibration sample, and also provide interesting physics. At the LHC, the ratio of production of top quark pairs and production of background to top quark pairs is more favorable than at the Tevatron, and clean samples of $t\bar{t}$ events can be selected, in particular in events with at least one electron or muon. ATLAS, for example, expects a signal as shown in Figure 6(right) in 100 pb$^{-1}$ of early data, even without b-tagging. Such a sample will be useful for the jet energy scale calibration, and calibration of b-tagging and $E_T$. With a working b-tagging, very clean $t\bar{t}$ samples can be selected.

The estimation of backgrounds in SUSY searches is performed as much as possible from data. As an example (Figure 6(center)) the important $Z (\rightarrow \nu\nu)$ plus jets background in the SUSY jets plus $E_T$ channel can be calibrated with a clean $Z (\rightarrow \mu\mu)$ plus jets sample, but also with $W (\rightarrow \mu\nu)$ plus jets samples. Many other control samples are under study.

6 Final comments

The first LHC data is eagerly awaited by a large community of experimentalists and theorists. There will be a strong pressure on the experiments to provide results early, and it is hard to prevent high-profile analyses to take place in a “glass box”. There will also be strong internal competition within ATLAS and CMS, and between ATLAS and CMS. It is important not to compromise on the quality of the results, after all the LHC will be operating in new territory. In this sense, the issue of blind analyses at the LHC, in order to prevent any bias, has come up, but has also raised internal discussions (at least in ATLAS) on the feasibility.

ATLAS and CMS are learning more and more from CDF and D0, not the least since many CDF and D0 physicists are joining ATLAS and CMS. There is still a lot to be learned from the Tevatron in terms of analysis techniques and background estimations from data, and on W and Z production. Certainly, CDF and D0 have shown that understanding the detector with data will be a major challenge.

References

4. A. Drozdetskiy, this conference
THE STRONGLY-INTERACTING LIGHT HIGGS

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We present the effective low-energy lagrangian arising from theories where the electroweak symmetry breaking is triggered by a light composite Higgs, which emerges from a strongly-interacting sector as a pseudo-Goldstone boson. This lagrangian proves to be useful for LHC and ILC phenomenology that includes the study of high-energy longitudinal vector boson scattering, strong double-Higgs production and anomalous Higgs couplings.

1 Introduction

The Standard Model (SM) of elementary particles, as we know it today, is not a complete theory. As it is well known, if we calculate the amplitude of the process $W_L W_L \rightarrow W_L W_L$ we find that it grows with the energy as $g^2 E^2 / M_W^2$ violating unitarity at energies around $4\pi v \sim 1$ TeV. What unitarize this amplitude at high energy? This is the first priority question to be addressed at the LHC.

An example of a possible UV-completion of the SM can be found in QCD. The pion amplitudes are unitarized by extra resonances arising from the strongly interacting SU(3)$_c$. Nevertheless, this Higgsless approach has to face the present electroweak precision test (EWPT) and, in its simple incarnation, technicolor models, it fails to pass them. The reason is that the new resonances responsible for unitarizing the SM amplitudes have masses at around 1 TeV and give large (tree-level) contribution to electroweak observables that have not been observed.

A second option arises from the Higgs mechanism. The presence of a scalar Higgs cures the SM amplitudes from the bad high-energy behaviour and, therefore, allow the SM to be extrapolated to very high energies. It is hard to believe that nature is not using such a simple mechanism to give us a UV completed theory of electroweak interactions. Nevertheless, naturalness criteria, stop us from considering the Higgs mechanism as the last ingredient to be incorporated to the SM at the electroweak scale. Why the Higgs mass, that determines the electroweak scale, is so small compare with, for example, the Planck scale? If we want to answer this question, we
must postulate new dynamics at the electroweak scale. For example, supersymmetry can give an explanation to the smallness of the electroweak scale. Nevertheless, no super-partner of the SM particles have been found at the present colliders making the supersymmetric solution less and less natural.

A third possibility that has received a big boost in the recent years is the composite Higgs idea. Similarly as pions in QCD, the Higgs is assumed to be a composite particle and therefore not suffering from any naturalness problem. The role of this Higgs is, however, less ambitious than its original motivation. Since the Higgs is composite, its coupling to the SM particles will be different from those of a point-like scalar. Therefore the SM amplitudes are only partly unitarized by the composite Higgs and extra resonances must be present in these models to completely unitarize the SM amplitudes. At this point, one could ask: What are we gaining? Well, the presence of a composite Higgs makes the SM viable up to a higher energy than in models without a Higgs, implying that the extra resonances can be heavier. Hence, their effects on the EW observables are smaller and these models can be able to safely pass all the EWPT.

In the first models of composite Higgs the Higgs appeared as a Pseudo-Goldstone boson (PGB) from a strong interacting theory, very similar to pions in QCD. These first proposals, however, lacked several ingredients. First, it did not incorporate a heavy top (since its mass was not known at that time). The authors of Ref. had to enlarge the SM gauge group to obtain EWSB. Second, the contribution to EW observables were not calculated due mainly, to the strong regime of the theory. Also flavor was also not successfully incorporated. Recently there has been various attempts to realize the composite Higgs idea avoiding the above problems. This includes the Little Higgs, Holographic Higgs and other variations.

Here we want to study the general properties and the phenomenology of scenarios in which a light Higgs is associated with strong dynamics at a higher scale, focusing on features that are quite independent of the particular model realization. We will refer to this scenario as to the Strongly-Interacting Light Higgs (SILH). Of course, in many specific models, the best experimental signals will be provided by direct production of new states, while here we concentrate on deviations from SM properties in Higgs and longitudinal gauge boson processes. Still, we believe that our model-independent approach is useful. The tests we propose here on Higgs and gauge-boson interactions will help, in case of new discoveries, to establish if the new particles indeed belong to a strongly-interacting sector ultimately responsible for electroweak symmetry breaking. If no new states are observed, or if the resonances are too broad to be identified, then our tests can be used to investigate whether the Higgs is weakly coupled or is an effective particle emerging from a strongly-interacting sector, whose discovery has been barely missed by direct searches at the LHC.

2 The structure of SILH

2.1 Definition of SILH

The structure of the theories we want to consider is the following. In addition to the vector bosons and fermions of the SM, there exists a new sector responsible for EW symmetry breaking, which is broadly characterized by two parameters, a coupling $g_p$ and a scale $m_p$ describing the mass of heavy physical states. Collectively indicating by $g_{SM}$ the SM gauge and Yukawa couplings (basically the weak gauge coupling and the top quark Yukawa), we assume $g_{SM} \lesssim g_p \lesssim 4\pi$. The upper bound on $g_p$ ensures that the loop expansion parameter $\sim (g_p/4\pi)^2$ is less than unity, while the limit $g_p \sim 4\pi$ corresponds to a maximally strongly-coupled theory in the spirit of naive dimensional analysis (NDA). Because of the first inequality, by a slight abuse of language, we shall refer to the new sector as “the strong sector”. The Higgs multiplet is assumed to belong to the strong sector. The SM vector bosons and fermions are weakly coupled to the strong
sector by means of the $SU(3) \times SU(2) \times U(1)_Y$ gauge coupling and by means of proto-Yukawa interactions, namely interactions that in the low-energy effective field theory will give rise to the SM Yukawas.

A second crucial assumption we are going to make is that in the limit $g_{SM} = 0$, $g_{\rho} \neq 0$ the Higgs doublet $H$ is an exact Goldstone boson, living in the $\mathcal{G}/\mathcal{H}$ coset space of a spontaneously broken symmetry of the strong sector. Two minimal possibilities in which the complex Higgs doublet spans the whole coset space are $SU(3)/SU(2) \times U(1)$ and the custodially symmetric $SO(5)/SO(4)$. The $\sigma$-model scale $f$ will be assumed to be related to $g_{\rho}$ and $m_{\rho}$ by the equation

$$m_{\rho} = g_{\rho} f.$$  \tag{1}

The gauging of $SU(2) \times U(1)_Y$ and the non-zero Yukawas explicitly break the Goldstone symmetry of the strong sector leading to terms in the (effective) action that are not invariant under the action of $\mathcal{G}$ on the coset space. In particular a mass term for the Higgs is generated at 1-loop. If the new dynamics is addressing the hierarchy problem, it should soften the sensitivity of the Higgs mass to short distances, that is to say below $1/m_{\rho}$. In interesting models, the Higgs mass parameter is thus expected to scale like $(\alpha_{SM}/4\pi)m_{\rho}^2$. Observation at the LHC of the new states with mass $m_{\rho}$ will be the key signature of the various realizations of SILH. Here, as stated before, we are interested in the model-independent effects, which could be visible in processes involving the Higgs boson and/or longitudinal gauge bosons, and which would unmistakably reveal new physics in the electroweak breaking sector.

### 2.2 The SILH Effective Lagrangian

Below $m_{\rho}$, the field content of these theories consists in the SM particles plus the Higgs. Deviations from the SM are encoded in the higher dimensional operators of the low-energy effective lagrangian. The dimension-6 operators involving the Higgs field can be separated in three parts, depending on the origin of the operators:

1. At tree-level $1/f^2$ order:

$$\frac{c_H}{2f^2} \partial^\mu \left( H^\dagger H \right) \partial_\mu \left( H^\dagger H \right) + \frac{c_T}{2f^2} \left( H^\dagger \overleftrightarrow{D^\mu} H \right) \left( H^\dagger \overleftrightarrow{D_\mu} H \right)$$

$$-c_{\rho}\frac{\lambda}{f^2} \left( H^\dagger H \right)^3 + \left( \frac{c_{\rho}g_{\rho}}{f^2} H^\dagger H f_L H f_R + \text{h.c.} \right).$$  \tag{2}

2. At tree-level $1/m_{\rho}^2$ order:

$$\frac{ig_{\rho} g}{2m_{\rho}^2} \left( H^\dagger \sigma^\mu \overleftrightarrow{D^\mu} H \right) (D^\nu W_{\mu\nu})^i + \frac{ic_{\rho} g'}{2m_{\rho}^2} \left( H^\dagger \overleftrightarrow{D^\mu} H \right) (\sigma^\mu B_{\mu\nu}).$$  \tag{3}

3. At one-loop order:

$$\frac{i c_{\rho} g f^2}{16\pi^2 f^2} (D^\mu H)^\dagger \sigma^\mu (D^\nu H) W^i_{\mu\nu} + \frac{i c_{\rho} g' f^2}{16\pi^2 f^2} (D^\mu H)^\dagger (D^\nu H) B_{\mu\nu}$$

$$+ \frac{c_{\rho} g'^2}{16\pi^2 f^2} \frac{f^2}{g_{\rho}^2} \left( H^\dagger H B_{\mu\nu} B_{\mu\nu} \right) + \frac{c_{\rho} g'^2}{16\pi^2 f^2} \frac{f^2}{g_{\rho}^2} \left( H^\dagger H C_{\mu\nu} C_{\mu\nu} \right).$$  \tag{4}

The coupling constants $c_i$ are pure numbers of order unity. For phenomenological applications, we have switched to a notation in which gauge fields are canonically normalized, and gauge couplings explicitly appear in covariant derivatives. Also, we recall the definition $H^\dagger \overleftrightarrow{D_\mu} H \equiv H^\dagger D_\mu H - (D_\mu H^\dagger) H$.
The first class of operators (those of order $1/f^2$) are the most sizable. Among them only the operator proportional to $c_T$ is constrained from the experimental data. Since it violates the custodial symmetry, it gives a contribution $\Delta \rho \equiv \tilde{T} = c_T \xi$, \begin{equation}
abla (\xi) = \frac{v^2}{f^2}, \quad v = \left(\sqrt{2} G_F\right)^{-1/2} = 246 \text{ GeV}.
\end{equation}
From the SM fit of electroweak data, we find $-1.1 \times 10^{-3} < c_T \xi < 1.3 \times 10^{-3}$ at 95% CL (letting also $S$ to vary one finds instead $-1.7 \times 10^{-3} < c_T \xi < 1.9 \times 10^{-3}$ at 95% CL). This strong limit on $c_T$ suggests that new physics relevant for electroweak breaking must be approximately custodial-invariant. In our Goldstone Higgs scenario this corresponds to assuming the coset $SO(5)/SO(4)$. When $g_{SM}$ is turned on, $c_T$ receives a model dependent contribution, which should be small enough to make the model acceptable. The rest of the coefficient $c_H$, $c_y$ and $c_\phi$ are practically unconstrained and their implications will be discussed in the next section.

A linear combination of the operators with coefficients $c_W$ and $c_B$ contributes to the $\tilde{S}$ parameter of electroweak precision data:
\begin{equation}
\tilde{S} = (c_W + c_B) \frac{m_W^2}{m_B^2},
\end{equation}
where $\tilde{S}$ is defined in ref. Using the SM fit of electroweak data, we obtain the bound $m_\rho > (c_W + c_B)^{1/2} 2.5$ TeV at 95% CL. (this bound corresponds to assuming a light Higgs and $\Delta \rho \equiv \tilde{T} = 0$; by relaxing this request the bound becomes $m_\rho > (c_W + c_B)^{1/2} 1.6$ TeV). In terms of the parameter $\xi$ defined in eq. (6), this bound becomes
\begin{equation}
\xi \lesssim \frac{1.5}{c_W + c_B} \left(\frac{g_\rho}{4\pi}\right)^2.
\end{equation}
As we show later, new effects in Higgs physics at the LHC appear only for sizable values of $\xi$. Then eq. (8) requires a rather large value of $g_\rho$, unless $c_W + c_B$ happens to be accidentally small.

The operators with coefficients $c_{HH}$ and $c_{BB}$ originate from the 1-loop action $\mathcal{L}^{(1)}$, under our assumption of minimal coupling for the classical action. Although they are $H^2 D^1$ terms, like $c_W$, $c_B$, they cannot be enhanced above their 1-loop size by the exchange of any spin 0 or 1 massive field. The operators proportional to $c_y$ and $c_\phi$ are suppressed by an extra power $(g_{SM}/g_\rho)^p$ with respect to those proportional to $c_{HH}$ and $c_{BB}$. Moreover, while $c_H$ and $c_y$ indirectly correct the physical Higgs coupling to gluons and quarks by $O(v^2/f^2)$ with respect to the SM, the direct contribution of $c_T$ and $c_\phi$ is of order $(v^2/f^2)(g_{SM}/g_\rho)^p$. Their effect is then important only in the weakly coupled limit $g_\rho \sim g_{SM}$. Notice that from the point of view of the Goldstone symmetry, $O_{BB}$ and $O_3$ are like a Higgs mass term with extra field strength insertions.

In the simplest models $m_H^2 \sim (g_{SM}^2/16\pi^2)m_\rho^2$. We have here assumed this simplest possibility, which accounts for the extra $g_{3}^2 m_3^2/g_\rho^2$ appearing in eq. (4). More precisely, for phenomenological purposes, we have chosen $g_{SM}^2$ as the coupling of the largest contribution in the corresponding SM loop, i.e., $g_{SM}^2 = g_\rho^2 (y^2 H)$ for the operator involving photons (gluons), respectively.

3 Phenomenology of SILH

In this section we analyze the effects of the SILH interactions and study how they can be tested at future colliders. Let us start by considering the new interaction terms involving the physical Higgs boson. For simplicity, we work in the unitary gauge and write the SILH effective Lagrangian in eqs. (2)-(4) only for the real Higgs field $h$ (shifted such that $h = 0$). We reabsorb
the contributions from \( c_f \) and \( c_b \) to the SM input parameters (fermion masses \( m_f \), Higgs mass \( m_H \), and vacuum expectation value \( v = 246 \text{ GeV} \)). Similarly, we redefine the gauge fields and the gauge coupling constants and we make the gauge kinetic terms canonical. In this way, the SILH effective Lagrangian is composed by the usual SM part, written in terms of the usual SM input parameters (physical masses and gauge couplings), by new Higgs interactions \( (\mathcal{L}_h) \), and new interactions involving only gauge bosons \( (\mathcal{L}_V) \) which, at leading order, are given by

\[
\mathcal{L}_h = \xi \left( \frac{c_H}{2} \left( 1 + \frac{1}{v} \right)^2 \partial \mu h \partial \mu h - c_t \frac{m_t^2}{v^2} \left( \frac{1}{2} v^2 + \frac{3}{2} t v \right) + c_b \frac{m_b^2}{v} f f + \frac{3}{2} h^2 \right) + \left( \frac{h}{v} + \frac{h^2}{2 v^2} \right) \left( \frac{g^2}{2 g_p^2} \left( \hat{c}_W W^\mu \partial \nu W^\nu - h.c. \right) + \frac{g^2}{2 g_p^2} Z^\mu \partial \nu Z^\nu \right) + \left( \frac{c_{HW}}{2} \right)^2 W^\mu W^\nu - \left( \frac{c_{HW} + \tan^2 \theta_W c_{HB} c_{H}}{4} \right) \left( \frac{g}{4 \pi} \right)^2 \left( \frac{c_{HW} + \tan^2 \theta_W c_{HB} c_{H}}{4} \right) W^\mu W^\nu \right)
\]

(9)

\[
\hat{c}_W = c_W + \left( \frac{g}{4 \pi} \right)^2 c_{HW}, \quad \hat{c}_B = c_B + \left( \frac{g}{4 \pi} \right)^2 c_{HB}, \quad \hat{c}_Z = c_Z + \frac{c_{EB} - c_{HW}}{4 \sin 2 \theta_W}
\]

(10)

\[
\mathcal{L}_V = \left( \frac{1}{2} \tan \frac{\theta_W}{2} \right) S W_\mu^{(3)} B_\mu - i g \cos \theta_W g_1 Z^\mu \left( W^\mu W^\nu - W^\nu W^\mu \right) - i g \left( \cos \theta_W \cos \kappa \kappa' A_\mu \right) W^\nu W^\nu
\]

(11)

In \( \mathcal{L}_V \) we have included only trilinear terms in gauge bosons and dropped the effects of \( O_{2W}, O_{2B}, O_{2Z} \). In \( \mathcal{L}_h \) we have kept only the first powers in the Higgs field \( h \) and the gauge fields. We have defined \( W^\mu = \partial_\mu W^\pm - \partial_\mu \partial W^\pm \) (and similarly for the \( Z_\mu \) and the photon \( A_\mu \)) and \( D_\mu = \partial_\mu \partial_\nu - \Box g_{\mu \nu} \). Notice that for on-shell gauge bosons \( D_\mu A^\mu = M_A^2 A_\mu \). Therefore \( \hat{c}_W \) and \( \hat{c}_B \) generate a Higgs coupling to gauge bosons which is proportional to mass, as in the SM, and do not generate any Higgs coupling to photons.

The new interactions in \( \mathcal{L}_h \) see eq. (9), modify the SM predictions for Higgs production and decay. At quadratic order in \( h \), the coefficient \( c_H \) generates an extra contribution to the Higgs kinetic term. This can be reabsorbed by redefining the Higgs field according to \( h \rightarrow h / \sqrt{1 - \xi c_H} \). The effect of \( c_{HW} \) is then to renormalize by a factor \( 1 - \xi c_H/2 \), the couplings of the canonical field \( h \) to all other fields. We can express the modified Higgs couplings in terms of the decay widths in units of the SM prediction, expressed in terms of physical pole masses,

\[
\Gamma (h \rightarrow f \bar{f})_{\text{SILH}} = \Gamma (h \rightarrow f \bar{f})_{\text{SM}} \left[ 1 - \xi \left( 2 c_y + c_H \right) \right]
\]

(13)

\[
\Gamma (h \rightarrow W^+ W^-)_{\text{SILH}} = \Gamma (h \rightarrow W^+ W^-)_{\text{SM}} \left[ 1 - \xi \left( c_H - \frac{g^2}{g_p^2} \hat{c}_W \right) \right]
\]

(14)

\[
\Gamma (h \rightarrow Z Z)_{\text{SILH}} = \Gamma (h \rightarrow Z Z)_{\text{SM}} \left[ 1 - \xi \left( c_H - \frac{g^2}{g_p^2} \hat{c}_Z \right) \right]
\]

(15)

\[
\Gamma (h \rightarrow gg)_{\text{SILH}} = \Gamma (h \rightarrow gg)_{\text{SM}} \left[ 1 - \xi \left( 2 c_y + c_H + \frac{4 g_y^2 c_y}{g^2_p} \right) \right]
\]

(16)
\begin{align}
\Gamma (h \to \gamma \gamma)_{\text{SILH}} &= \Gamma (h \to \gamma \gamma)_{\text{SM}} \left[ 1 - \xi \Re \left( \frac{2c_y + c_H}{1 + I_\gamma/J_\gamma} + \frac{c_H - \frac{4}{g_\rho} \xi c_W}{1 + I_\gamma/J_\gamma} + \frac{4c_y c_\gamma}{I_\gamma + J_\gamma} \right) \right] \\
\Gamma (h \to \gamma Z)_{\text{SILH}} &= \Gamma (h \to \gamma Z)_{\text{SM}} \left[ 1 - \xi \Re \left( \frac{2c_y + c_H}{1 + J_\gamma/I_\gamma} + \frac{c_H - \frac{4}{g_\rho} \xi c_W}{1 + J_\gamma/I_\gamma} + \frac{4c_y c_\gamma}{I_\gamma + J_\gamma} \right) \right].
\end{align}

Here we have neglected in $\Gamma (h \to W^+W^-, ZZ)_{\text{SILH}}$ the subleading effects from $c_{H W}$ and $c_{H B}$, which are parametrically smaller than a SM one-loop contribution. The loop functions $I$ and $J$ are given in Ref.\(^5\).

The leading effects on Higgs physics, relative to the SM, come from the three coefficients $c_H$, $c_y$, $c_{\gamma Z}$, although $c_{\gamma Z}$ has less phenomenological relevance since it affects only the decay $h \to \gamma Z$. Therefore, we believe that an important experimental task to understand the nature of the Higgs boson will be the extraction of $c_H$ and $c_y$ from precise measurements of the Higgs production rate ($\sigma_h$) and branching ratios ($BR_h$). The contribution from $c_H$ is universal for all Higgs couplings and therefore it does not affect the Higgs branching ratios, but only the total decay width and the production cross section. The measure of the Higgs decay width at the LHC is very difficult and it can be reasonably done only for rather heavy Higgs bosons, well above the two gauge boson threshold, while the spirit of our analysis is to consider the Higgs as a pseudo-Goldstone boson, and therefore relatively light. However, for a light Higgs, LHC experiments can measure the product $\sigma_h \times BR_h$ in many different channels: production through gluon, gauge-boson fusion, and top-strahlung; decay into $b$, $\tau$, $\gamma$ and (virtual) weak gauge bosons. At the LHC with about 300 fb\(^{-1}\), it is possible to measure Higgs production rate times branching ratio in the various channels with 20–40 % precision\(^6\), although a determination of the $b$ coupling is quite challenging\(^8\). This will translate into a sensitivity on $|c_H \xi|$ and $|c_y \xi|$ up to 0.2–0.4.

In fig. 1, we show our prediction for the relative deviation from the SM expectation in the main channels for Higgs discovery at the LHC, in the case $c_H \xi = 1/4$ and $c_y / c_H = 1$ (as in the Holographic Higgs). For $c_y / c_H = 0$, the deviation is universal in every production channel and is given by $\Delta (\sigma \ BR) / (\sigma \ BR) = -c_H \xi$.

Cleaner experimental information can be extracted from ratios between the rates of processes...
with the same Higgs production mechanism, but different decay modes. In measurements of these ratios of decay rates, many systematic uncertainties drop out. Our leading-order \( (g_p \gg g_{SM}) \) prediction is that \( \Delta \Gamma(h \to ZZ)/\Gamma(h \to W^+W^-) = 0, \Delta \Gamma(h \to f\bar{f})/\Gamma(h \to W^+W^-) = -2\xi_{ctv} \), \( \Delta \Gamma(h \to \gamma\gamma)/\Gamma(h \to W^+W^-) = -2\xi_{ctv}(1 + J_\gamma/I_\gamma)^{-1} \). However, the Higgs coupling determinations at the LHC will still be limited by statistics, and therefore they can benefit from a luminosity upgrading, like the SLHC. At a linear collider, like the ILC, precisions on \( \sigma_t \propto BR_h \) can reach the percent level \( ^9 \), providing a very sensitive probe on the new-physics scale. Moreover, a linear collider can test the existence of \( c_h \), since the triple Higgs coupling can be measured with an accuracy of about 10\% for \( \sqrt{s} = 500 \) GeV and an integrated luminosity of 1000 fb\(^{-1}\)\(^{10} \).

Deviations from the SM predictions of Higgs production and decay rates, could be a hint towards models with strong dynamics, especially if no new light particles are discovered at the LHC. However, they do not unambiguously imply the existence of a new strong interaction. The most characteristic signals of a SILH have to be found in the very high-energy regime. Indeed, a peculiarity of SILH is that, in spite of the light Higgs, longitudinal gauge-boson scattering amplitudes grow with energy and the corresponding interaction becomes strong, eventually violating tree-level unitarity at the cutoff scale. Indeed, the extra Higgs kinetic term proportional to \( c_H \xi \) in eq. (9) prevents Higgs exchange diagrams from accomplishing the exact cancellation, present in the SM, of the terms growing with energy in the amplitudes. Therefore, although the Higgs is light, we obtain strong \( WW \) scattering at high energies.

From the operator \( O_{H} \equiv \mathcal{G}(H^+H^+)\partial_{\mu}(H^+H^+) \) in eq. (2), using the equivalence theorem \(^{11} \), it is easy to derive the following high-energy limit of the scattering amplitudes for longitudinal gauge bosons

\[
A(Z_L^0 Z_L^- \to W_L^+ W_L^-) = A(W_L^+ W_L^- \to Z_L^0 Z_L^0) = \frac{c_{ft} m_H^2}{f^2}, \\
A(W_L^\pm W_L^- \to W_L^\pm W_L^\pm) = \frac{c_H(s + t)}{f^2}, \quad A(W_L^+ W_L^- \to W_L^+ W_L^-) = \frac{c_H(s + t)}{f^2},
\]

(19) \hspace{2cm} (20) \hspace{2cm} (21)

This result is correct to leading order in \( s/f^2 \), and to all orders in \( \xi \) in the limit \( g_{SM} = 0 \), when the \( \sigma \)-model is exact. The absence of corrections in \( \xi \) follows from the non-linear symmetry of the \( \sigma \)-model, corresponding to the action of the generator \( T_h \), associated with the neutral Higgs, under which \( \nu \) shifts. Therefore we expect that corrections can arise only at \( O(s/m_H^2) \). The growth with energy of the amplitudes in eqs. (19)-(21) is strictly valid only up to the maximum energy of our effective theory, namely \( m_{W} \). The behaviour above \( m_{W} \) depends on the specific model realization. In the case of the Little Higgs, we expect that the amplitudes continue to grow with \( s \) up to the cut-off scale \( \Lambda \). In 5D models, like the Holographic Goldstone, the growth of the elastic amplitude is softened by KK exchange, but the inelastic channel dominate and strong coupling is reached at a scale \( \sim 4\pi m_W/g_p \). Notice that the result in eqs. (19)-(21) is exactly proportional to the scattering amplitudes obtained in a Higgsless SM \(^{11} \). Therefore, in theories with a SILH, the cross section at the LHC for producing longitudinal gauge bosons with large invariant masses can be written as

\[
\sigma(pp \to V_L V_L^\pm X)_{ct} = (c_H \xi)^2 \sigma(pp \to V_L V_L^\pm X)_{H},
\]

where \( \sigma(pp \to V_L V_L^\pm X)_{H} \) is the cross section in the SM without Higgs, at the leading order in \( s/(4\pi v)^2 \). With 200 fb\(^{-1} \) of integrated luminosity, it should be possible to identify the signal of a Higgsless SM with about 30-50\% accuracy \(^{12,13} \), i.e., to a sensitivity up to \( c_H \xi \approx 0.5-0.7 \).

In the SILH framework, the Higgs is viewed as a pseudo-Goldstone boson and therefore its properties are directly related to those of the exact (eaten) Goldstones, corresponding to
the longitudinal gauge bosons. Thus, a generic prediction of SILH is that the strong gauge boson scattering is accompanied by strong production of Higgs pairs. Indeed we find that, as a consequence of the $O(4)$ symmetry of the $H$ multiplet, the amplitudes for Higgs pair-production grow with the center-of-mass energy as eq. (19),

$$A (Z_L^0 Z_L^0 \to hh) = A (W_L^+ W_L^- \to hh) = \frac{e_H s}{f^2}.$$  

(23)

Notice that scattering amplitudes involving longitudinal gauge bosons and a single Higgs vanish. This is a consequence of the $Z_4^X$ parity embedded in the $O(4)$ symmetry of the operator $O_H$, under which each Goldstone change sign. Non-vanishing amplitudes necessarily involve an even number of each species of Goldstones.

Using eqs. (19), (20) and (23), we can relate the Higgs pair production rate at the LHC to the longitudinal gauge boson cross sections

$$2\sigma_{\Delta M} (pp \to hh X)_{\epsilon_H} = \sigma_{\Delta M} (pp \to W_L^+ W_L^- X)_{\epsilon_H} + \frac{1}{6} \left( 9 - \text{tanh}^4 \frac{\delta}{2} \right) \sigma_{\Delta M} (pp \to Z_L^0 Z_L^0 X)_{\epsilon_H}.$$  

(24)

Here all cross sections $\sigma_{\Delta M}$ are computed with a cut on the pseudorapidity separation between the two final-state particles (a boost-invariant quantity) of $|\Delta \eta| < \delta$, and with a cut on the two-particle invariant mass $\hat{s} > M^2$. The sum rule in eq. (24) is a characteristic of SILH. However, the signal from Higgs-pair production at the LHC is not so prominent. It was suggested that, for a light Higgs, this process is best studied in the channel $bb\gamma\gamma$\(^{14}\), but the small branching ratio of $h \to \gamma\gamma$ makes the SILH rate unobservable. However, in SILH, one can take advantage of the growth of the cross section with energy. Although we do not perform here a detailed study, it may be possible that, with sufficient luminosity, the signal of $bbbb$ with high invariant masses could be distinguished from the SM background.

acknowledgments

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Beyond the Standard Model Higgs in ATLAS and CMS

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Theories beyond the Standard Model, such as Supersymmetry, the Littlest Higgs Model and models with Extra Dimensions predict the existence of Higgs bosons of different kinds such as CP-even and CP-odd neutral Higgs, singly charged and doubly charged Higgs, CP-violating Higgs, radions and invisibly decaying Higgs. A number of recent analyses of simulated searches for such Higgs bosons at LHC are reviewed, stating in each case conditions under which conditions a 5σ discovery can be made.

1 Introduction

However successful the Standard Model (SM) is to describe with high precision almost all experimental observations made at the energies so far accessible at high energy colliders, it is at the same time plagued by instabilities due to quadratically divergent loop corrections at higher energies. Furthermore, the Standard Model has no candidate particle to explain Dark Matter, the existence of which is based on firm astronomical and cosmological observations. It can furthermore not give a satisfactory explanation of the total dominance of matter over antimatter in Universe.

Different schemes beyond the Standard Model have been proposed to solve several or all of these problems. These extensions of the Standard Model in general have a Higgs sector that is more complex than that of the Standard Model. Among the most popular schemes are:

- Supersymmetry, which introduces heavy superpartners to all SM particles. The superpartners have loop corrections that are precisely of the same magnitude as those of the SM particles but of opposite sign and thereby the divergences are cancelled to all orders. As the Standard Model instabilities set in already at a scale of 1 TeV, the SUSY symmetry braking scale needs to be of the same order.

- The Little Higgs model, which introduces a set of heavier vector bosons and top-antitop quarks that provide a limited cancellation and push the divergences up to order 10 TeV.
• Models with extra dimensions, which has the SM interactions confined to the four ordinary dimensions whereas gravity occupies also the extra dimensions.

2 Higgs Bosons in Supersymmetry

With Supersymmetry more than 100 extra model parameters are introduced. For practical applications the model, therefore, has to be constrained, like it is in the constrained Minimal Supersymmetric Model (MSSM), which is the minimal supersymmetric extension of the Standard Model with 7 free parameters. In MSSM two complex Higgs field doublets are introduced which is the minimal Higgs structure needed to keep the theory anomaly free and to give mass to the fermions. MSSM predicts the existence of five physical Higgs bosons, three neutral Higgs, of which two, h and H, are CP-even and one, A, is CP-odd, and two charged Higgs bosons, H⁺ and H⁻. In order to further constrain the number of free parameters four benchmark scenarios have been proposed¹ for which 5 of the 7 parameters are given fixed values. The predictions of the theory can then be displayed as function of the two remaining free parameters, often chosen as tanβ, the ratio of the vacuum expectation value of the two Higgs doublets, and MA, the mass of the CP-odd Higgs boson. The values of the other 5 parameters in the four proposed benchmark scenarios are shown in Table 1.

<table>
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<tr>
<th>Name</th>
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<td>-1100</td>
<td>500</td>
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Table 1: Values of the SUSY parameters in the four benchmark scenarios

Below are discussed a few recent specific MSSM results obtained from simulation studies with ATLAS and CMS.

3 Associate production of H/A and b̅b, with the H/A decaying to μ⁺μ⁻

For large tanβ, the process gg → b̅bΦ, where Φ is one of the MSSM neutral Higgs h, H or A, represents a dominant Higgs production process and provides the best measurements of the heavy Higgs H and A. Different Φ decay channels have been investigated: Φ → μ⁺μ⁻, Φ → τ⁺τ⁻ with the τ decaying to e, μ and jets and Φ → b̅b. The τ⁺τ⁻ final state has the largest Φ discovery reach whereas the μ⁺μ⁻ final state provides the best mass resolution, comparable in magnitude to the natural width of the Higgs, which is a sensitive measure of tanβ. The presence of two b quarks is used to suppress background from μ⁺μ⁻ Drell-Yan production. A simulation analysis of the Abb associate production and A → μ⁺μ⁻ decay has been made by CMS². Figure 1 (left) shows the peak of a 150 GeV mass A Higgs boson above the background and the square points in Figure 1 (right) show the width of this A peak and its measurement error as function of tanβ. The round points in Figure 1 (right) represent a calculation in the MSSM of the A width as function of tanβ and the triangular points represent the addition to the A width, resulting from the degeneracy of the A and H masses at low tanβ values, which leads to an apparent widening of the observed A peak.
4 Search for charged Higgs

4.1 Charged Higgs decay to $\tau\nu_{\tau}$

It has been shown that the charged Higgs decay channel $H^{+} \to \tau^{+}\nu_{\tau}$ can be used to search for $H^{+}$ up to Higgs masses just below the top mass, irrespectively of the value of $\tan\beta$. For such mass values $H^{+}$ is produced through top decays $t \to bH^{+}$. For $H^{+}$ masses greater than the top mass the dominant production process is through gluon fusion. The Higgs decay $H^{+} \to \tau^{+}\nu_{\tau}$ remains the best search channel up to large $H^{+}$ masses. The $H^{+} \to t\bar{b}$ decay is less efficient as it suffer from a large QCD background.

A discovery reach analysis has been made in ATLAS using a high $p_T$ jet and large missing transverse momentum at large angle as distinctive feature. In the transition region around the top mass the two processes $gg \to t\bar{b}H^{+}$ and $gb \to tH^{+}$ both contribute. They partially represent the same process, which leads to double counting. This problem has recently been solved with the MC program MadHiggs in which the doubly counted events are subtracted. The results of the ATLAS analysis show that with 30 fb$^{-1}$ integrated luminosity a 5σ discovery can be made of a 170 GeV mass $H^{+}$ for $\tan\beta$ values above 20 and of a 600 GeV mass $H^{+}$ for $\tan\beta$ values above 40.

4.2 Charged Higgs decay to SUSY particles

LEP has excluded $\tan\beta$ values below $\approx 4$ for all $M_A$ values. The previously mentioned $H^{+} \to \tau^{+}\nu_{\tau}$ search analysis covers the region down to $\tan\beta \approx 20$. For intermediate $\tan\beta$ values the discovery of $H^{+}$, decaying to Standard Model particles only, is made difficult by a minimum in the $t\bar{b}H^{+}$ coupling of around $\tan\beta \approx 7.5$. An analysis of the discovery potential for a $H^{+}$ decaying in the chargino-neutralino channel has recently been made in ATLAS with the goal to cover the intermediate $\tan\beta$ region. For this analysis was chosen a MSSM parameter point for which the $H^{+}$ decay branching ratio to a chargino and a neutralino is maximized. The $H^{+}$ production and decay diagrams are shown in Figure 2 (left). The parameter values of the chosen MSSM point are: $M_2 = 210$ GeV, $\mu = 135$ GeV, $m_{l_R} = 110$ GeV, $m_{l_K} = 210$ GeV, $m_{\tilde{g}} = 800$ GeV and $m_{\tilde{q}} = 1000$ GeV. The signature used for the selection of the signal is: a) three hard leptons, b) substantial missing transverse energy due to the neutralinos and the neutrino and c) three hard jets from the hadronic top decay. The resulting 5σ $H^{+}$ discovery region in the $\tan\beta-M_A$ parameter plane is shown in Figure 2 (right). An earlier result obtained for CMS is also shown in this figure. As can be seen, according to this analysis, $H^{+}$ of mass value from 250 to 400 GeV can be discovered also in the previously uncovered region of $\tan\beta$ values from 5 to 20.
5 MSSM with CP violation

The massive terms appearing in the soft SUSY braking terms can be complex with non-trivial CP-violating phases. One of the consequences of such CP violation is that the three neutral MSSM Higgs $h$, $H$ and $A$ mix into states $H_1$, $H_2$ and $H_3$ that have no definite CP parity. Such states could not be detected at LEP since CP-violating Higgs do not couple to the $Z$ boson. The discovery potential for the specific CP-violating MSSM benchmark parameter point called CPX\(^7\) has been investigated in ATLAS\(^8\) by reinterpretation of studies made within the SM and MSSM. The resulting $5\sigma$ discovery regions in the $\tan\beta$ vs $m_{H^\pm}$ plane for $H_1$, $H^+$ and $H_2$ and/or $H_3$, respectively, are shown in Figure 3 (left) for an integrated luminosity of 300 fb\(^{-1}\). $H_1$ can be discovered in most of the free $\tan\beta$ vs $m_{H^\pm}$ plane shown in Figure 3 and together with $H^+$, $H_2$ and $H_3$ the whole plane is covered except for a tiny area around $\tan\beta = 5$ and $m_{H^\pm} = 145$ GeV, which is shown in the blow-up around this parameter point in Figure 3 (right).

6 Higgs searches in non-SUSY models

6.1 Search for the radion

In the Randall-Sundrum model\(^9\) there are two four-dimensional hyper surfaces or branes located at the boundary of the fifth dimension. The Standard Model fields are on the visible brane and all mass terms, which are at order Planck mass, are rescaled to order TeV by an exponential factor. The fluctuations in the metric in the fifth dimension are described by a radion scalar field $\Phi$ which has a vacuum expectation value $\Lambda_\Phi$. The radion has mass $m_\Phi$ and mixes with the Higgs boson $h$ with a mixing parameter $\xi$.

Radions can be produced from gluon fusion and decay to two $h$ bosons; $gg \to \Phi \to hh$. The final states of interest are $\gamma\gamma b\overline{b}$ and $\tau^+\tau^- b\overline{b}$. A search analysis has been made by CMS\(^10\) assuming $m_h = 125$ GeV, $m_\Phi = 300$ GeV and $\Lambda_\Phi = 1000$ GeV. Figure 4 (left) shows the distribution of effective mass of the $\gamma\gamma$ system, displaying a peak at the $h$ mass, which is
comparatively very sharp due to the high γ energy resolution in the CMS electromagnetic calorimetry. Figure 4 (centre) shows the effective mass of the γγbb system with a broad peak at the Φ mass. The larger width in this case is caused by the comparatively much lower mass resolution in the reconstruction of the h decaying to bb. In Figure 4 (right) is shown the resulting 5σ discovery contour in the A_Φ vs ξ parameter plane, the inner and outer boundaries corresponding to the result with and without systematic errors, respectively.

![Graphs showing effective mass and discovery contour](image)

Figure 4: left: The γγ effective mass in the gg → radion Φ → hh → γγbb process, centre: The γγbb effective mass in the gg → radion Φ → hh → γγbb process, right: The 5σ discovery contour of the radion with (inner boundary) and without (outer ditto) systematic errors

6.2 Search for the doubly charged Higgs of the Littlest Higgs model

In the Little Higgs model extra heavy gauge bosons W' and Z' and a vector like heavy quark pair T, T̅ are introduced to stabilize the one-loop quadratically divergent radiative corrections. The Littlest Higgs model\(^{11}\), the minimal version of this model, predicts the existence of a doubly charged Higgs Δ⁺⁺ which can be Drell-Yann produced in pairs. The Δ⁺⁺ can decay into like sign leptons.

A study has been made by CMS in which the pair-produced doubly charged Higgs decay into four muons\(^{12}\). Background processes are t̅t → W⁺W⁻bb → 4μ, Zbb → 4μ, ZZ → 4μ and ZZ → 2μ 2τ. The upper Δ⁺⁺ mass limit for a 5σ discovery with 10 fb⁻¹ integrated luminosity was found to be 650 GeV and the upper mass limit for a 2σ exclusion 780 GeV.

6.3 Invisible Higgs

The neutral Higgs boson may in certain models decay into undetectable particles. Such invisibly decaying Higgs appear in MSSM (decay into the lightest neutralinos and gravitinos) and also, e.g., in models with large extra dimensions or in the extended SM with extra singlet scalars discussed at this meeting by Joachim van der Bij\(^{13}\). Higgs decaying into undetectable particles can only be triggered and reconstructed at LHC from particles produced in association with the Higgs.

The discovery potential for an invisibly decaying Higgs produced in association with a Z, with the Z decaying into leptons, has been investigated in ATLAS\(^{14}\). The dominant irreducible background is ZZ → llνν. This is a counting experiment (no mass peak reconstruction possible) and the detection of a signal relies on the accuracy of the absolute normalization, which is found to be of the order 4% at 100 fb⁻¹. Higgs of seven different Higgs mass values between 120 and 400 GeV were investigated at two integrated luminosities, 30 and 100 fb⁻¹. The results show that among these combinations it is only the one of 120 GeV mass and 100 fb⁻¹ integrated luminosity that would give a significance for the Higgs signal exceeding 5σ. Earlier investigations based on the invisibly decaying Higgs produced in the Vector Boson Fusion production mode with a final state containing an invisibly decaying Higgs and two hard quark jets\(^{15}\), have shown that a 5σ signal can, for this production mode, be obtained over most of the mass range investigated with 100 fb⁻¹ integrated luminosity. However, for this production mode, the discrimination on the trigger level represents a major problem.
7 Concluding remarks

Although the first signs at LHC of physics beyond the Standard Model may not necessarily come from the detection of non-Standard-Model Higgs bosons, the discovery of such bosons, and the measurement of their properties, will under all circumstances be of crucial importance for understanding the New Physics. It is therefore essential to find out how Higgs bosons of the various theories beyond the Standard Model can be discovered and their properties measured at the LHC. Even if the analyses using simulated ATLAS and CMS data are now in quite an advanced state with many interesting new results, it is urgent to advance these analyses further to enable a rapid interpretation of the signals from the new, unexpected and maybe surprising physics that will hopefully appear at LHC.

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Symmetry Breaking in Six Dimensional Flux Compactification Scenarios

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Motivated by the electroweak hierarchy problem, we consider theories with two extra dimensions in which the four-dimensional scalar fields are components of gauge boson in full space, namely the Gauge-Higgs unification framework. We briefly explain the basics features of "flux compactification", i.e. compactification in presence of a background (magnetic) flux. In particular we recall how chirality and symmetry breaking can be obtained in this context. More in details, we find and catalogue all possible degenerate zero-energy stable configurations in the case of trivial or non-trivial 't Hooft flux, for a SU(N) gauge theory on a torus. We describe the residual symmetries of each vacua and the four-dimensional effective spectrum in terms of continuous and discrete parameters, respectively.

1 Introduction

All available data agree in indicating that the mass of the Higgs boson is of the order of the Electroweak scale, \( v \sim O(200) \) GeV. Such a mass is unnaturally light if there is new physics beyond the Standard Model (SM) to which the Higgs boson is sensitive. In fact the Higgs mass parameter is not protected by any symmetry and thus gets corrections which are quadratically dependent on possible higher scales, like the unification scale \( M_{\text{GUT}} \) or, ultimately, the Planck scale \( M_{\text{Pl}} \).

Three different mechanisms have been devised in order to eliminate the quadratic sensitivity of the Higgs mass to the cutoff scale:

- **Supersymmetry**: bosonic and fermionic contributions to the quadratic divergences cancel each other in such a way that the Higgs mass remains affected only by a logarithmic sensitivity to the cutoff scale;

- **Technicolor** and **Little Higgs**: the Higgs is a Goldstone boson of a global custodial symmetry that it is only softly (spontaneously) broken;

- **Gauge-Higgs Unification**: the Higgs is a component of a higher dimensional gauge multiplet. The lightness of its mass is protected by the gauge symmetry itself.

Independently of the precise nature assumed for the Higgs field, all these proposals require, in one way or another, the appearance of new physics at about the TeV scale. While the first two approaches are being intensely studied, in practice they tend to be afflicted by rather severe fine-tuning requirements when confronted with present experimental data. Here, instead, we concentrate on the last and less explored possibility: Gauge-Higgs unification. The idea is that a single higher dimensional gauge field gives rise to all the four-dimensional (4D) bosonic degrees of freedom: the gauge bosons, from the ordinary space-time components and the scalar bosons.
(and the Higgs fields among them) from the extra ones. The essential point concerning the solution of the hierarchy problem is that, although the higher dimensional gauge symmetry is globally broken by the compactification procedure, however it always remains locally unbroken. Any local (sensitive to the UV physics) mass term for the scalars is then forbidden by the gauge symmetry and the Higgs mass would then only have a non-local and UV finite origin.

This idea has been widely investigated in the context of five- and six-dimensional orbifold compactification\(^2\). From the field theory point of view, a different and less explored possibility is to recover the idea of Gauge-Higgs unification in the context of flux compactification: compactification of the extra space-like dimensions on a manifold in which there exist a (gauge) background with a non-trivial field strength, compatible with Scherk-Schwarz periodicity conditions\(^3\). We'll review in the following the basic idea of five- and six-dimensional SS compactifications.

2 Scherk-Schwarz mechanism in five- and six-dimensional compactifications

Let's consider a \(U(N)\) gauge theory on a \((4 + d)\)-dimensional space-time\(^a\) where the extra dimensions are compactified on an orthogonal \(d\)-dimensional torus \(T^d\). To completely define a field theory on a torus one has to specify the periodicity conditions; that is, to describe how the fields transform under the fundamental shifts \(y \rightarrow y + l_a\), with \(l_a\) being the lengths of the non-contractible cycles of the torus. Let's denote with \(T_a\) the embeddings of these shifts in the fundamental representation of \(U(N)\). If we want to preserve four-dimensional Poincaré invariance, the twists \(T_a\) must depend only on the extra-dimensional coordinates. The most general periodicity conditions for the gauge field \(A_M\) and for a generic field \(\Phi\) in the fundamental representation of \(U(N)\), read respectively:

\[
A_M(x, y + l_a) = T_a(y) \left[ A_M(x, y) + \frac{i}{g} \partial_M \right] T^*_a(y), \tag{1}
\]

\[
\Phi(x, y + l_a) = T_a(y) \Phi(x, y). \tag{2}
\]

These equations are derived from the fact that while individual (gauge or matter) fields may not be single-valued on the torus, any physical scalar quantity, like the Lagrangian, must be. The periodicity conditions in Eqs. (1,2) are usually referred as Scherk-Schwarz (SS) boundary conditions\(^3\). Let's describe in more details the five- and six-dimensional compactification procedure in presence of general SS boundary conditions.

2.1 Scherk-Schwarz mechanism in five-dimensions

In the case of a five-dimensional theory compactified on a circle \(S^1\) one has to define a single twist matrix \(T(y)\). No restrictions have to be imposed on \(T\) except that it belongs to the \(U(N)\) gauge group\(^5\). The \(U(N)\) twist matrix can be, locally, decomposed as the product of an element \(e^{e(x)} \in U(1)\) and an element \(V(y) \in SU(N)\) as follows:

\[
T(y) = e^{e(x)} V(y). \tag{3}
\]

It is always possible to choose a gauge, called the symmetric gauge\(^5\), in which the \(SU(N)\) vacuum configurations are trivial and the twist matrix is constant \(V^{sym} = V\) and can be parametrized as:

\[
V = e^{2\pi i (\alpha \cdot H)}, \quad \alpha \cdot H \equiv \sum_{j=1}^{N-1} \alpha^j H_j, \tag{4}
\]

\(^a\)Throughout the paper, with \(x\) and \(y\) we will denote the four-ordinary and \(d\)-extra coordinates, respectively. Latin upper case indices \(M, N\) will run over all the extra-dimensional space, whereas greek and latin lower case indices \(\mu, \nu\) and \(a, b\) will run over the four ordinary and the extra-dimensions, respectively.

\(^b\)The case of external automorphisms is not considered here. See for example\(^4\)
where $H_j$ are the $(N-1)$ generators of the Cartan subalgebra of $SU(N)$ and $\alpha_j$ are $(N-1)$ real continuous parameters $0 \leq \alpha_j < 1$. These parameters are non-integrable phases, which arise only in a topologically non-trivial space and cannot be gauged-away. When all the $\alpha_j$ are vanishing the periodicity conditions are trivial and consequently the initial symmetry is unbroken. If, instead, some of the $\alpha_j$ are non-vanishing, then symmetry breaking can occur. This mechanism is known as the Scherk-Schwarz mechanism\(^3\).

In order to give an explicit expression for the gauge masses, one introduces the Cartan-Weyl basis for the $SU(N)$ generators. In addition to the Cartan subalgebra generators, $H_j$, one defines $N(N-1)$ non-diagonal generators, $E_r$, such that the following commutation relations are satisfied:

$$[H_j, H_k] = 0 \ , \ [H_j, E_r] = q_r^j E_r .$$

(5)

In this basis, the twist $V$ acts in a diagonal way, that is

$$V H_j V^\dagger = H_j \quad , \quad V E_r V^\dagger = e^{2\pi i (\alpha \cdot q_r)} E_r ,$$

(6)

and the four-dimensional mass spectrum reads simply:

$$m_{l_h}^2 = \frac{4\pi^2}{l^2} (n + \alpha \cdot q_h)^2 \quad , \quad n \in \mathbb{Z} .$$

(7)

For field components associated to a generator belonging to the Cartan subalgebra, $H_j$, one has $q_j = (0, ..., 0)$ and the spectrum reduce to the ordinary Kaluza-Klein (KK) one. For field components associated to the non-diagonal generators, $E_r$, one has, instead, $q_r \neq (0, ..., 0)$ and the mass spectrum is consequently shifted by a factor proportional to the non-integrable phases $\alpha^j$. When all the $\alpha^j \neq 0$, then only the gauge field components associated to the generators of the Cartan subalgebra are massless. Therefore, the symmetry breaking induced by the twists, $V$, does not lower the rank of $SU(N)$. The maximal symmetry breaking pattern that can be achieved for an $U(N)$ symmetry group is given by:

$$U(N) \sim U(1) \times SU(N) \rightarrow U(1) \times U(1)^{N-1} = U(1)^N .$$

(8)

Scherk-Schwarz symmetry breaking mechanism can be used to break both global (flavour symmetries, supersymmetry) or local symmetries. In the case of gauge symmetry breaking the SS phase, $\alpha$, can be interpret as the vev of the extra-dimensional component of the gauge fields, $\langle A_5 \rangle$. At classical level the scalar potential is flat and consequently the phases $\alpha^j$ are undetermined. Their values must be dynamically determined at the quantum level\(^{10,5}\) minimizing the one-loop effective potential. If, at the minimum, any of the $\alpha$ is non-vanishing then the gauge symmetry is spontaneously broken. This dynamical and spontaneous symmetry breaking mechanism is conventionally known as the Hosotani mechanism. At the same time, the extra-dimensional component of the gauge field, $A_5$, is a scalar field that can be identified with the Higgs field and that acquires a finite mass term. The non-local nature of this symmetry breaking protects the theory from ultraviolet divergences and makes it a promising candidate mechanism to break the electroweak symmetry and to provide an Higgs field free from quadratic divergences.

2.2 Scherk-Schwarz mechanism in six-dimensions

In the case of a six-dimensional theory compactified on a torus $T^2$, one can introduce a different twist, $T_a(y)$, along each of the two independent cycles. The twists cannot be chosen arbitrarily but they have to satisfy the following $U(N)$ 't Hooft consistency condition\(^6,7\):

$$T_1(y + l_2) T_2(y) = T_2(y + l_1) T_1(y) .$$

(9)
This condition is obtained imposing that (for any fields included in the theory) the value of the field at the final point \((y_1 + l_1, y_2 + l_2)\), starting from the initial point \((y_1, y_2)\) has to be independent on the followed paths.

The \(U(N)\) twist matrices can be, locally, decomposed as the product of an element \(e^{\text{tw}_a(y)} \in U(1)\) and an element \(V_a(y) \in SU(N)\) as follows:

\[
T_a(y) = e^{\text{tw}_a(y)} V_a(y).
\]

Using this parametrization, the consistency conditions in Eq. (9) can be splitted into the \(SU(N)\) and \(U(1)\) part, respectively:

\[
V_1(y + l_2) V_2(y) = e^{2\pi i \frac{m}{N}} V_2(y + l_1) V_1(y)
\]
\[
\Delta_2 v_1(y) - \Delta_1 v_2(y) = 2\pi \frac{m}{N},
\]

with \(\Delta_a v_b(y) = v_b(y + l_a) - v_b(y)\). The \(SU(N)\) consistency condition, Eq. (11), tells us that the twists, \(V_a\) must commute, on the fundamental plaquette, modulo a phase factor belonging to the center of \(SU(N)\). The integer \(m = 0, 1, \ldots, (N - 1)\) (modulo \(N\)) is a gauge invariant quantity called the non-abelian \(\text{'t Hooft flux}\). Furthermore, Eq. (12) tells us that it must coincide with the value of a quantized abelian magnetic flux living on the torus or, in other words, with the first Chern class of \(U(N)\) on \(T^2\).

It is well known, that the presence of a stable magnetic background, associated with the abelian subgroup \(U(1) \subset U(N)\) and living only on the two extra dimensions, can induce chirality\(^8\) in four-dimensions. A non-vanishing value of the \(\text{'t Hooft flux}\) \(m\) is indeed necessary for having four-dimensional chiral matter fields. A general description of fermions and chirality in the context of 6D \(U(N)\) theories compactified on a two-dimensional torus can be found in\(^9\).

From the other side, from Eqs. (11,12) it appears evident that the presence of the quantized abelian magnetic flux deeply affects the non-abelian subgroup \(SU(N) \subset U(N)\), giving rise to a non-trivial \(\text{'t Hooft non-abelian flux}\). While the symmetry breaking pattern for a \(SU(N)\) theory in presence of trivial non-abelian \(\text{'t Hooft flux}\) \((m = 0)\) is well-known in the literature\(^{10,11}\), the field theory and phenomenological analysis of the non-trivial \((m \neq 0)\)\( \text{'t Hooft flux}\) has been explored only recently, in\(^{11,12}\). Here, it has been shown that exists a gauge, denominated, as the five dimensional case, the symmetric gauge, in which the \(SU(N)\) twists can always be chosen as constant, i.e. \(V_a^{\text{sym}} = V_a\), with \(V_a\) constant matrices satisfying the \(SU(N)\) \(\text{'t Hooft} consistency conditions:

\[
V_1 V_2 = e^{2\pi i \frac{m}{N}} V_2 V_1.
\]

In the symmetric gauge, the \(SU(N)\) vacuum configurations are trivial and therefore the residual symmetries of each classical vacua are those associated to the \(SU(N)\) generators which commute simultaneously with \(V_1\) and \(V_2\). The number of classical vacua and the pattern of symmetry breaking depend on the values of \(m\) and they will be analyzed in the following section.

3 \(SU(N)\) Symmetry Breaking: trivial vs non-trivial \(\text{'t Hooft flux}\)

The main purpose this section is to find and classify all possible vacua and to describe the residual symmetries for an effective four-dimensional theory obtained from a \(SU(N)\) gauge theory on a six-dimensional space-time where the two extra dimensions are compactified on a torus, for both the cases of trivial and non-trivial \(\text{'t Hooft} non-abelian flux.\)
3.1 Trivial ’t Hooft flux: $m = 0$

In the $m = 0$ case, Eq. (13) tells us that the two $V_a$ matrices commute and consequently can be parametrized as:

$$V_a = e^{2\pi i (\alpha_a \cdot H)} , \quad \alpha_a \cdot H = \sum_{j=1}^{N-1} \alpha_a^j H_j$$

with $H_j$ the $(N - 1)$ generators of the Cartan subalgebra of $SU(N)$. The periodicity conditions, and consequently the classical vacua, are now characterized by $2(N - 1)$ real continuous parameters, $0 \leq \alpha_a^j < 1$. As in the five-dimensional case these parameters are non-integrable phases, which arise only in a topologically non-trivial space and cannot be gauged-away. When all the $\alpha_a^j$ are vanishing the initial symmetry is unbroken. At classical level $\alpha_a^j$ are undetermined. Their values must be dynamically determined at the quantum level where the rank-preserving Hosotani symmetry breaking mechanism can occur.

The mass spectrum of the four-dimensional gauge and scalar components of the 6D gauge field follows straightforwardly the five-dimensional discussion. In the Cartan-Weyl basis Eq. 5, the twists $V_a$ act in a diagonal way, that is

$$V_a H_j V^\dagger_a = H_j , \quad V_a E_r V^\dagger_a = e^{2\pi i (\alpha_a \cdot \phi_r)} E_r ,$$

and the four-dimensional mass spectrum for gauge/scalar fields reads:

$$m^2_{(k)} = 4\pi^2 \sum_{a=1}^{2} (n_a + \alpha_a \cdot q_k)^2 \frac{1}{l_a^2} , \quad n_a \in \mathbb{Z} .$$

This is the same kind of spectrum seen previously in the five-dimensional case. For gauge (scalar) field components associated to a generator belonging to the Cartan subalgebra, $H_j$, the spectrum reduce to the ordinary Kaluza-Klein (KK) one. For gauge (scalar) field components associated to the non-diagonal generators, $E_r$, the mass spectrum is consequently shifted by a factor proportional to the non-integrable phases $\alpha_a^j$. Therefore, the symmetry breaking induced by the commuting twists, $V_a$, does not lower the rank of $SU(N)$.

One can easily generalize these results to the $U(N)$ case adding an extra diagonal generator, $H_0 = 1_N/\sqrt{2N}$. Obviously $H_0$ commute with all the twists $V_a$ and consequently $A^0_M$ always remains unbroken. The maximal symmetry breaking pattern that can be achieved in the $m = 0$ case, for an $U(N)$ gauge theory is given by:

$$U(N) \sim U(1) \times SU(N) \rightarrow U(1) \times U(1)^{N-1} = U(1)^N .$$

This symmetry breaking mechanism is exactly the same Hosotani mechanism one is used to in a five-dimensional framework.

3.2 Non-trivial ’t Hooft flux: $m \neq 0$

In the $m \neq 0$ case, the twists $V_a$ don’t commute between themselves and so necessarily they induce a rank-reducing symmetry breaking. The most general solution of the consistency relation Eq. (13) can be parametrized as follows:

$$V_1 = \omega_1 P^{s_1} Q^{t_1} , \quad V_2 = \omega_2 P^{s_2} Q^{t_2} .$$

$s_{a}, t_{a}$ are integers parameters taking values between $0, ..., (N - 1)$ (modulo $N$) and satisfying the following constraint:

$$s_1 t_2 - s_2 t_1 = m/K \equiv \bar{m} .$$
\( P \) and \( Q \) are \( SU(N) \) constant matrices given by
\[
P \equiv P_{\tilde{N}} \otimes \mathbb{1}_K , \quad Q \equiv Q_{\tilde{N}} \otimes \mathbb{1}_K
\]
where \( K = \text{g.c.d.}(m,N) \) and \( \tilde{N} = N/K \). \( P_{\tilde{N}} \) and \( Q_{\tilde{N}} \) are \( \tilde{N} \times \tilde{N} \) matrices defined as
\[
\begin{align*}
(P_{\tilde{N}})_{kj} &= e^{i\pi \tilde{N}^{-1}} \delta_{k,j-1} , & k, j = 1, 2, \ldots, \tilde{N}, \\
(Q_{\tilde{N}})_{kj} &= e^{i\pi \tilde{N}^{-1} (k+1)} e^{i\pi \tilde{N}^{-1}} \delta_{kj},
\end{align*}
\]
and satisfying the conditions
\[
P_{\tilde{N}} Q_{\tilde{N}} = e^{-2i\pi \tilde{N}} Q_{\tilde{N}} P_{\tilde{N}} , \quad (P_{\tilde{N}})_{\tilde{N}} = (Q_{\tilde{N}})_{\tilde{N}} = e^{i\pi \tilde{N}^{-1}}.
\]
When \( K = 1 \), then \( \tilde{N} = N \) and \( P, Q \) reduce to the usual elementary twist matrices defined by 't Hooft in \(^{6}\).

The matrices \( \omega_\alpha \) are constant elements of \( SU(K) \subset SU(N) \). They commute between themselves and with \( P \) and \( Q \). Therefore \( \omega_\alpha \) can be parametrized in terms of generators \( H_j \) belonging to the Cartan subalgebra of \( SU(K) \):
\[
\omega_\alpha = e^{i\pi (\alpha_\alpha H)} , \quad \alpha_\alpha \cdot H = \sum_{\rho=1}^{K-1} \alpha_\rho H_\rho
\]
Here \( \alpha_\rho \) are \( 2(K-1) \) real continuous parameters, \( 0 \leq \alpha_\rho < 1 \). As in the \( m = 0 \) case, they are non-integrable phases and their values must be dynamically determined at the quantum level producing a dynamical and spontaneous symmetry breaking.

The \( m \neq 0 \) four-dimensional mass spectrum is easily obtained using the following basis \(^{12}\) for the \( SU(N) \) generators
\[
\tau_{(\rho,\sigma)}(\Delta, k_{\Delta}) = \begin{cases} 
\text{if } \begin{cases} \rho = \sigma \\
\Delta = k_{\Delta} = 0
\end{cases} & \Rightarrow \left( \sum_{i=1}^{\rho} \lambda_{(i,i)}^{\rho} - \rho \lambda_{(\rho+1,\rho+1)}^{\rho} \right) \otimes \mathbb{1}_{\tilde{N}} \\
\text{else} & \Rightarrow \lambda_{(\rho,\sigma)}^{\rho} \otimes \tau_{\tilde{N}}(\Delta, k_{\Delta})
\end{cases}
\]
where \( \Delta, k_{\Delta} \) are integers assuming values between 0, \ldots, \( (\tilde{N} - 1) \) while the indices \( \rho, \sigma \) take values between 1, \ldots, \( K \), excluding the case \( \Delta = k_{\Delta} = 0, \rho = \sigma \) in which \( \rho \) takes values between 1, \ldots, \( (K - 1) \). The matrices \( \lambda_{(\rho,\sigma)}^{\rho} \) and \( \tau_{\tilde{N}} \) are \( K \times K \) and \( \tilde{N} \times \tilde{N} \) matrices, respectively, defined as:
\[
\lambda_{(\rho,\sigma)}^{\rho}(k_{\Delta}) = \delta_{\rho,\rho} \delta_{\sigma,\sigma} \\
\tau_{\tilde{N}}(\Delta, k_{\Delta}) = \sum_{n=1}^{\tilde{N}} e^{2i\pi \Delta n / \tilde{N}} \lambda_{(n,n+\Delta)}^{\tilde{N}}.
\]
The definition of \( \lambda_{(n,n+\Delta)}^{\tilde{N}} \) comes straightforwardly.

In this basis, the \( SU(K) \) generators that commute with \( P \) and \( Q \) are simply given by \( \tau_{(\rho,\sigma)}(0,0) \). In particular, the generators belonging to the Cartan subalgebra of \( SU(K) \) are given by \( H_\rho = \tau_{(\rho,\rho)}(0,0) \). The following commutation relations are satisfied:
\[
[\tau_{(\rho,\rho)}(0,0), \tau_{(\rho,\rho)}(0,0)] = 0 , \quad [\tau_{(\rho,\rho)}(0,0), \tau_{(\rho,\sigma)}(\Delta, k_{\Delta})] = \delta_{\rho,\rho} \delta_{\sigma,\sigma} \tau_{(\rho,\sigma)}(\Delta, k_{\Delta}).
\]
The action of the twists $V_a$ on this basis is given by

$$V_a \tau_{(\rho,\sigma)}(\Delta, k_\Delta) V_a^\dagger = \frac{2\pi i (n_a \Delta + t_a k_\Delta) + 2\pi i (\alpha_0 q_{(\rho,\sigma)})}{\tau_{(\rho,\sigma)}(\Delta, k_\Delta)}, \quad (26)$$

and the four-dimensional mass spectrum takes the following form:

$$m_{(\rho,\sigma)}^2(\Delta, k_\Delta) = 4\pi^2 \sum_{a=1}^{2} \left( n_a + \frac{1}{N} (s_a \Delta + t_a k_\Delta) + \alpha_0 \cdot q_{(\rho,\sigma)} \right)^2 \frac{1}{t_a^2}, \quad (27)$$

with $n_a \in \mathbb{Z}$. Therefore, beside the usual KK mass term, there are other two additional contributions. The first one, quantized in terms of $1/N$, is a consequence of the non-trivial commutation rule of Eq. (22) between $P$ and $Q$ that induces the $SU(N) \to SU(\mathcal{K})$ symmetry breaking. Since $s_a, t_a$ cannot be simultaneously zero, the spectrum described by Eq. (27) always exhibits some (tree-level) degree of symmetry breaking. Given a set of $s_a, t_a$ and for all the $\alpha_0 = 0$ (that is $\omega_a = 1$), only the gauge bosons components associated to $\tau_{(\rho,\sigma)}(0,0)$, the generators of $SU(\mathcal{K})$, admit zero modes. This is an explicit breaking. The second contribution to the gauge mass comes from the $\omega_a$ and it depends on the continuous parameters $\alpha_0^\rho$. For $\mathcal{K} > 1$ and all the non-integrable phases $\alpha_0^\rho \neq 0$, the only massless modes correspond to the gauge bosons associated to the Cartan subalgebra of $SU(\mathcal{K})$, i.e. $\tau_{(\rho,\rho)}(0,0)$. The symmetry breaking pattern induced by the $\omega_a$ produce a Hosotani symmetry breaking that does not lower the rank of $SU(\mathcal{K})$.

The maximal symmetry breaking pattern that can be achieved for an $U(N)$ gauge theory with matter fields in the fundamental is, in the $m \neq 0$ case, given by:

$$U(N) \sim U(1) \times SU(N) \to U(1) \times U(1)^{\mathcal{K}-1} = U(1)^{\mathcal{K}}. \quad (28)$$

Obviously, when $\mathcal{K} = 1$ the $SU(N)$ subgroup is completely broken, the only unbroken symmetry being the $U(1) \in U(N)$. This symmetry breaking pattern has no analogous in 5-dimensional frameworks and it’s peculiar of higher dimensional models where (topological) fluxes can appear.

As a final comment on the spectrum, notice that in both the cases of trivial and non-trivial ’t Hooft flux, the classical effective four-dimensional spectrum depends on the gauge indices but it does not depend on the Lorentz ones. This implies that at the classical level the 4D scalar fields $A_\mu$, arising from the extra-components of a six-dimensional gauge fields, are expected to be degenerate with the 4D gauge fields $A_\mu$ with the same gauge quantum numbers. This degeneracy is always removed at the quantum level.

4 Conclusions

In this paper we have analyzed possible symmetry breaking mechanism in the context of Gauge-Higgs unification scenario. The introduction of general five-dimensional SS boundary conditions can drive a 4D gauge symmetry breaking through the dynamical mechanism, conventionally known as Hosotani mechanism. One-loop contributions to the scalar sector can shift the minimum of the effective potential and generate a non-vanishing vev for the Higgs field. This symmetry breaking is spontaneous and rank preserving.

In six dimensions, SS boundary conditions have to satisfy a consistency condition. We discussed in details the $U(N)$ case where a novel ingredient appears: the non-abelian ’t Hooft flux. This flux is a topological quantity intimately connected with the $U(1)$ (quantized) magnetic flux. In the case of trivial $(m = 0)$ ’t Hooft flux the gauge symmetry breaking obtained though SS boundary condition is the usual rank preserving Hosotani mechanism. In the case of non-trivial $(m \neq 0)$ ’t Hooft flux one can have, instead, two simultaneous symmetry breaking mechanism. A explicit, rank reducing, symmetry breaking associated to the non-commutativity of the twists leading to the $SU(N) \to SU(\mathcal{K})$ breaking. On top of that, for $\mathcal{K} > 1$, the residual symmetry group can be further reduced through a spontaneous, rank preserving, Hosotani mechanism.
The simultaneous presence of rank preserving and rank reducing symmetry breaking mechanism makes the non-trivial 't Hooft flux case particularly interesting from a model building point of view.

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References

A GRAVITY DUAL OF SINGLE-SECTOR SUPERSYMMETRY BREAKING

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We present a five-dimensional gravitational dual of "single-sector" models of supersymmetry breaking which are models that contain no messenger sector and naturally explain the scale of supersymmetry breaking and the fermion mass hierarchy. Inspired by flux-background solutions of type IIB supergravity, supersymmetry is broken by a metric background that deviates from AdS_5 in the infrared. The first and second generation sfermions directly feel the supersymmetry breaking and obtain masses of order 10 TeV, while the gauginos and third generation sfermions are elementary states that obtain soft masses of order 1 TeV at the loop level via direct gauge mediation. This particle spectrum leads to distinctive signatures at the LHC, similar to the usual gauge mediation with a neutralino NLSP that decays promptly to a gravitino LSP.

1 Introduction

Recently, gauge/gravity duality ideas based on the AdS/CFT correspondence in type IIB string theory have been used to give a four-dimensional (4D) holographic description for models in a warped extra dimension. This has led to the remarkable result that strongly coupled 4D gauge dynamics can be modeled with a five-dimensional (5D), weakly coupled gravitational theory. In this approach, classical field theory computations are able to capture the dominant effects of the strongly coupled 4D theory.

Warped extra dimension models have been used to break supersymmetry, where the warp factor is used to generate a low SUSY-breaking scale which is then identified as a dynamical supersymmetry breaking (DSB) scale. In these previous models boundary conditions were used to break supersymmetry. In the present work, we pursue this idea of relating the warp factor with a dynamically generated scale in the context of realistic, strongly coupled 4D SUSY gauge theories, softly broken by the effects of DSB. A simple 5D gravitational dual will be described
that will allow previously "incalculable" particle mass spectra to be calculated.

Supersymmetry will be broken by considering an effective 5D model that is motivated from a ten-dimensional (10D) type IIB supergravity solution, which is obtained by perturbing the well-known Klebanov-Strassler supersymmetric background using techniques developed in Ref. The effective deformed nonsupersymmetric 5D background metric is obtained from a dimensional reduction of the 10D metric, and the resulting 5D geometry will be parametrized as

$$ds^2 = A^2(z)(-dt^2 + dx_5^2 + dz^2),$$

where the warp factor is

$$A^2(z) = \frac{1}{(kz)^2} \left[ 1 - \epsilon \left( \frac{z}{z_1} \right)^4 \right],$$

$k$ is the $AdS$ curvature scale, and $z_0 \leq z \leq z_1$ with $z_0$, $z_1$ the positions of the (ultraviolet) UV and (infrared) IR branes respectively. The parameter $\epsilon$ is related to variables in the original 10D solution (see Ref.), although for our phenomenological purposes we only need assume it to be an arbitrary but small, positive parameter. The $\epsilon \to 0$ limit is just a slice of $AdS_5$, which is the 5D background setup used in the Randall-Sundrum model.

2 A 5D Gravity Model

The MSSM fields are assumed to propagate in the bulk with metric (1). In the supersymmetric limit ($\epsilon \to 0$) these 5D fields propagate in a slice of $AdS$ and satisfy nontrivial boundary conditions. Upon compactification to four dimensions, the massless zero modes of the Kaluza-Klein towers are identified with the 4D MSSM fields. Since the warp factor is used to set the scale of supersymmetry breaking, the Higgs fields need not be localized on the IR brane and in fact they are assumed to be confined on the UV brane where their masses are protected by supersymmetry.

2.1 Fermion masses

Consider the SM fermions, where each fermion is embedded into its own 5D Dirac spinor. The zero mode profile for each fermion $i$ is given by $f_i(z) \propto z^{\frac{3}{2} - c_i}$, where the exponent depends on a bulk mass parameter $c_i$. For $c_i > 1/2$ ($c_i < 1/2$) the zero mode is localized near the UV (IR) brane. The wavefunction overlap of the fermion zero modes with the UV-confined Higgs fields $z_0 = 1/k$, leads to the 4D Yukawa couplings

$$Y_{\psi} = Y_{\psi}^{5D} k \sqrt{\frac{1/2 - c_L}{(kz_1)^{1-2c_L} - 1}} \sqrt{\frac{1/2 + c_R}{(kz_1)^{1+2c_R} - 1}},$$

where $Y_{\psi}^{5D}$ is a 5D Yukawa coupling. This expression is used to solve for the $c$ parameters using the values of the 4D Yukawa couplings and assuming $10^{-3} \leq Y_{\psi}^{5D} k \leq 1$. The results are listed in Table 1. Indeed it is seen from these values that the lighter generations are closer to the IR brane while the third generation is UV-localized. Since each SM fermion is contained in a chiral supermultiplet, the corresponding scalar superpartner will be localized at the same place in the supersymmetric limit. In the deformed case, the scalar localization is qualitatively unchanged. This is because the profile is only modified in the IR, where the deformation is noticeable.

Note that it is necessary to allow a small hierarchy in the 5D Yukawa couplings, in order to avoid FCNC’s from the squarks. Essentially, the $c$’s must be degenerate among first and second generation quarks in order for the corresponding squarks to be degenerate. This also helps to avoid naturalness constraints from hypercharge Fayet-Illiopoulos D-terms.
Table 1: Standard Model $\mathcal{N}=8$ running fermion masses at the scale $m_Z$ with the corresponding $c$ values and 5D Yukawa couplings (in units of $k$) for the case of UV Higgses and $\tan\beta = 10$.

<table>
<thead>
<tr>
<th>Particle</th>
<th>$\tilde{m}(m_Z)$</th>
<th>$c_L$</th>
<th>$c_R$</th>
<th>$\lambda^{SU}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$e$</td>
<td>0.503 MeV</td>
<td>0.350</td>
<td>0.350</td>
<td>1</td>
</tr>
<tr>
<td>$\mu$</td>
<td>103.9 MeV</td>
<td>0.467</td>
<td>0.467</td>
<td>1</td>
</tr>
<tr>
<td>$\tau$</td>
<td>1.75 GeV</td>
<td>0.601</td>
<td>0.601</td>
<td>1</td>
</tr>
<tr>
<td>$d$</td>
<td>3.9 MeV</td>
<td>0.456</td>
<td>0.456</td>
<td>0.059</td>
</tr>
<tr>
<td>$s$</td>
<td>67.6 MeV</td>
<td>0.456</td>
<td>0.456</td>
<td>1</td>
</tr>
<tr>
<td>$b$</td>
<td>2.9 GeV</td>
<td>0.69</td>
<td>0.648</td>
<td>1</td>
</tr>
<tr>
<td>$u$</td>
<td>1.7 MeV</td>
<td>0.456</td>
<td>0.456</td>
<td>0.0025</td>
</tr>
<tr>
<td>$c$</td>
<td>0.58 GeV</td>
<td>0.456</td>
<td>0.456</td>
<td>0.849</td>
</tr>
<tr>
<td>$t$</td>
<td>166 GeV</td>
<td>0.69</td>
<td>5.341</td>
<td>1</td>
</tr>
</tbody>
</table>

2.2 Scalar masses

The scalar superpartner masses are related to the fermion mass spectrum. The zero mode profile of a bulk scalar field is given at leading order by $f_s(z) \propto z^{b_s-1}$, where the exponent depends on a mass parameter $b_s$ of the 5D model. By supersymmetry,

$$b_s = \frac{3}{2} - c_s,$$

which explicitly shows that once the SM fermion localization is set by $c_s$, the localization of the scalar zero mode is then fixed. The values $b_s < 1$ ($c_s > 1/2$) correspond to a UV-localized mode, whereas $b_s > 1$ ($c_s < 1/2$) is IR-localized. Clearly it is the IR-localized scalar modes that are sensitive to the SUSY-breaking background because the deformation is only appreciable near the IR brane.

In the supersymmetry-breaking background (1) the scalar zero modes will obtain a mass. The scalar mass squared as a function of the localization parameter $b$ is given by

$$\tilde{m}^2 = \epsilon \frac{(1-b)(b+10) (kz_1)^{1+b} - (kz_1)^{1-b}}{(kz_1)^{1+b} - (kz_1)^{1-b}} k^2 + \mathcal{O}(\epsilon^2).$$

This expression simplifies in the limit $kz_1 \gg 1$. For $b > 1$ the scalar mass simply becomes

$$\tilde{m} \approx \sqrt{\epsilon(b-1)(b+10)z_1^{-1}},$$

while for $0 < b < 1$ we have the approximation:

$$\tilde{m} \approx \sqrt{\epsilon(1-b)(b+10)(kz_1)^{b-1}z_1^{-1}}.$$  

Thus we see that for an IR-localized field ($b > 1$) the scalar mass becomes of order the IR scale $z_1^{-1}$, while for $b \ll 1$ and $kz_1 \sim 10^{13}$ the scalar mass is much less than a GeV.

From Eq. (4), the values of $b_s$ are determined by the fermion spectrum of Table 1. We then apply (5) to obtain the squark and slepton mass spectrum. The AdS curvature scale is set by requiring $m_3^2 \simeq M_5^2/k$ where $M_5$ is the 5D Planck scale. Choosing $k \sim 0.1 M_5$ requires $k \simeq 10^{-3/2} m_P = 7.7 \times 10^{10} \text{ GeV}$. Consequently the model parameters are set to

$$\pi k R = 28.42, \quad \epsilon = 0.05, \quad \tan\beta = 10, \quad z_0 = k^{-1}, \quad z_1 = (k e^{-x R})^{-1} = (35 \text{ TeV})^{-1}.$$  

The first two generations of squarks and sleptons obtain masses of order $1/10$ to $1/20$ the Kaluza-Klein mass scale,

$$m_{KK} = \pi z_1^{-1} = 110 \text{ TeV}.$$
but the third generation masses are much smaller. As expected since the third generation fermions are near the UV brane in order to have a large overlap with the Higgs, the corresponding supersymmetry-breaking masses are phenomenologically unacceptable. However by considering the dual 4D theory we will show that there is a gauge-mediated contribution that gives rise to acceptable third generation squark and slepton masses.

2.3 Relation to 4D single-sector models

Interestingly, the holographic 4D dual of the 5D gravity model is remarkably similar to models constructed purely in four dimensions. In particular, the authors of the “single-sector” models\textsuperscript{11,12} consider a class of theories in which DSB can be argued convincingly, and in which the first two generations of the MSSM arise as composite states (P\hat{U}) of a strongly coupled gauge theory. The fields \hat{U} acquire large F-terms, so that the composites (P\hat{U}) feel the SUSY-breaking directly. The first and second generation scalars get large masses, whereas the fermion composites remain massless due to chiral symmetries. Since the \hat{U} fields also carry Standard Model charges, they communicate SUSY-breaking to the rest of the MSSM through gauge mediation. The scalar masses for the first and second generation composite scalars (P\hat{U}) are given by \(m_\phi^2 \sim F^2/M^2\), where the messenger scale \(M\) is the scale of the strong internal dynamics, corresponding to the Kaluza-Klein scale (9) in the gravitational dual.

Taking into account the parameters chosen in (8), the messenger scale is thus

\[ M = 110 \text{ TeV}. \] (10)

We will assume \(F \approx M\), as is common in theories where the messengers couple strongly to the DSB sector. We also require a large enough \(F/M\) in order to have a viable spectrum, and this too leads to \(F \approx M\). In particular, we choose \(F/M = 90 \text{ TeV}\).

The other ingredient that is needed to compute the effects of gauge mediation is the number of messengers, \(N_m\). In practice we set \(N_m = 2\), since this gives rise to an attractive LHC phenomenology and satisfies the experimental constraints that will be discussed below. We note that the \(B\mu\) term and \(A\) terms are generated radiatively, with the boundary condition that they vanish at the messenger scale (10). This is fairly constraining and significantly influences the model parameters. In particular, the model is adjusted so as to obtain viable electroweak symmetry breaking and the lightest Higgs mass.

2.4 The particle mass spectrum

In Table 2 we show the complete soft mass spectrum using the two-loop RGE code Softsusy\textsuperscript{10}, for the values of the parameters given in (8), and \(\mu < 0\) for the Higgsino mass parameter. Boundary conditions are imposed at the messenger scale (10), and the bulk soft masses calculated from (5) are added in quadrature to the gauge mediation masses at that scale. Softsusy automates a self-consistent determination of the thresholds for the superpartner spectrum, taking into account one- and two-loop effects. The gravitino mass is obtained from the standard formula \(m_{3/2} = F/(\sqrt{3} m_F) = 2.35 \text{ eV}\).

Note that these are only the masses for the lightest modes, which are zero modes in the AdS\textsubscript{5} limit. The Kaluza-Klein modes are at the \(O(100) \text{ TeV}\) scale. The heavy first and second generation scalar masses arising from the bulk 5D calculation represent nonperturbative masses in the 4D dual theory that are difficult to calculate directly in the strongly coupled gauge theory.

2.5 LHC signal

The gravitino is the LSP, which means that in the single sector model the lightest neutralino, \(\tilde{\chi}_1^0\), is the NLSP. Because the messenger scale is relatively low, the decay length of \(\tilde{\chi}_1^0\) is less
Table 2: Particle mass spectrum of the example single-sector model described in the text.

<table>
<thead>
<tr>
<th>(\tilde{e}_L, \tilde{e}_R, \tilde{\nu}_L)</th>
<th>10160, 10150, 10160 GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\tilde{\mu}_L, \tilde{\mu}_R, \tilde{\nu}_L)</td>
<td>5145, 5130, 5145 GeV</td>
</tr>
<tr>
<td>(d_L, d_R, \tilde{d}_L, \tilde{d}_R)</td>
<td>5905, 5885, 5970, 5890 GeV</td>
</tr>
<tr>
<td>(\tilde{s}_L, \tilde{s}_R, \tilde{c}_L, \tilde{c}_R)</td>
<td>5905, 5885, 5970, 5890 GeV</td>
</tr>
<tr>
<td>(\tilde{g})</td>
<td>1615 GeV</td>
</tr>
<tr>
<td>(\tilde{b}_1, \tilde{b}_2, \tilde{t}_1, \tilde{t}_2)</td>
<td>1354, 1369, 1253, 1369 GeV</td>
</tr>
<tr>
<td>(\tilde{\tau}_1, \tilde{\tau}<em>2, \tilde{\nu}</em>{eL})</td>
<td>511, 630, 633 GeV</td>
</tr>
<tr>
<td>(\tilde{\chi}_1^{\pm}, \tilde{\chi}_2^{\pm})</td>
<td>478, 593 GeV</td>
</tr>
<tr>
<td>(\tilde{\chi}_1^0, \tilde{\chi}_2^0, \tilde{\chi}_3^0, \tilde{\chi}_4^0)</td>
<td>288, 480, 511, 598 GeV</td>
</tr>
<tr>
<td>(h^0, A^0, H^0, H^\pm)</td>
<td>115, 646, 646, 651 GeV</td>
</tr>
<tr>
<td>(\tilde{G})</td>
<td>2.35 eV</td>
</tr>
</tbody>
</table>

than 1 mm. This leads to the signal \(pp \to 2\gamma + E_T\) (two hard photons and missing transverse energy) at the LHC. The study of the diphoton signal can be performed using PYTHIA (version 64.08).\(^{13}\)

Most of the background in the diphoton channel can be removed by cuts on \(E_T\) and \(p_T\), since Standard Model diphoton events are predominantly of low \(p_T\) and \(E_T\). Hence we impose the following kinematic cuts to reduce background:

\[
p_T,\gamma \geq 40 \text{ GeV}, \quad E_T \geq 60 \text{ GeV}. \tag{11}\]

The results are shown in Fig. 1. It can be seen that backgrounds (dashed) are orders of magnitude smaller than the signal (solid). With 1-10 fb\(^{-1}\) of data, virtually no background events occur. The simple \(p_T,\gamma\) and \(E_T\) cuts (11) suffice to remove virtually all SM backgrounds for the diphoton plus missing energy signal. Discovery of the example model within the first 10 fb\(^{-1}\) of well-understood data is a certainty, and would occur during the first few years of the LHC experiment.

3 Conclusion

We have presented a 5D dual gravity model of 4D single-sector supersymmetry breaking models. These models naturally explain the scale of supersymmetry breaking and the fermion mass hierarchy without invoking a messenger sector. They lead to a distinctive particle spectrum consisting of heavy (\(O(10^2\text{eV})\)) first and second generation squark and slepton masses. The remaining sparticles are lighter (\(O(\text{TeV})\)) so that at low energies only the gluinos, charginos, neutralinos and third generation squarks and sleptons will be accessible at the LHC. The LSP is the gravitino. The most striking signal at the LHC is from diphotons and missing energy, which will be easily detectable after 1-10 fb\(^{-1}\) of "well-understood" data is accumulated.

The dual 4D interpretation of our model is that the first two generations of fermions and bosons would be composite states of some strongly coupled gauge theory ("superglue") that is responsible for both the scale of supersymmetry breaking via dimensional transmutation and the fermion mass hierarchy via large anomalous dimensions for fermionic operators in the gauge theory. The remaining particles are elementary fields that couple weakly to the composite supersymmetry breaking sector. This holographic interpretation is qualitatively identical to single-sector models that were explicitly constructed in four dimensions\(^{11,12}\). Our 5D model not only has a calculational advantage over 4D strongly coupled gauge theories, where at best only naive dimensional analysis estimates are possible, but also uses the AdS/CFT correspondence to identify the ratio of the Planck scale to the scale of supersymmetry breaking with the warp factor and the fermion mass hierarchy as arising from wavefunction overlap in the bulk.
Figure 1: Comparison of the single-sector diphoton signal (solid) to background (dashed) where cuts (11) have been made, removing virtually all the background.

Acknowledgments

It is a pleasure to thank Maxime Gabella and Joel Giedt for collaboration on the work presented here. This work was supported in part by a Department of Energy grant, No. DE-FG02-94ER40823 at the University of Minnesota, and by an award from the Research Corporation.

References

Electroweak Interactions and Unified Theories

III - Beyond SM, Extradim

Top Physics
SINGLE TOP RESULTS FROM CDF

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The CDF Collaboration has analyzed 955 pb$^{-1}$ of CDF II data collected between March 2002 and February 2006 to search for electroweak single top quark production at the Tevatron. We employ three different analysis techniques to search for a single top signal: multivariate likelihood functions; neural networks; the matrix element analysis technique. The sensitivities to a single top signal at the rate predicted by the Standard Model are 2.1 $\sigma$, 2.0 $\sigma$ and 2.5 $\sigma$, respectively. The first two analyses observe a deficit of single top-like events and set upper limits on the production cross section. The matrix element analysis observes a 2.3 $\sigma$ single top excess and measures a combined t-channel and s-channel cross section of $2.7^{+1.2}_{-1.4}$ pb. Using the same dataset, we have searched for non-Standard Model production of single top quarks through a heavy $W'$ boson resonance. No evidence for a signal is observed. We exclude at the 95% C.L. $W'$ boson production with masses of 760 GeV/$c^2$ (790 GeV/$c^2$) in case the right handed neutrino is smaller (larger) than the mass of the $W'$ boson.

1 Introduction

In 1.96 TeV proton anti-proton collisions at the Tevatron, top quarks are predominantly produced in pairs via the strong force. In addition, the Standard Model predicts single top quarks to be produced through an electroweak t- and s-channel exchange of a virtual $W$ boson (Figure 1). The production cross sections have been calculated at Next-to-Leading-Order (NLO). For a top quark mass of 175 GeV/$c^2$ the results are $1.98 \pm 0.25$ pb and $0.88 \pm 0.11$ pb for the t-channel and s-channel process respectively$^1$. The combined cross section is about 40% of the top antitop pair production cross section. The precise measurement of the production cross section allows the direct extraction of the Cabibbo-Kobayashi-Maskawa matrix element $|V_{tb}|$ and offers a source of almost 100% polarized top quarks$^2$. Moreover, the search for single top also probes exotic models beyond the Standard Model. New physics, like flavor-changing neutral currents or heavy $W'$ bosons, could alter the observed production rate$^3$. Finally, single top processes
result in the same final state as the Standard Model Higgs boson process $WH \rightarrow Wbb$, which is one of the most promising low mass Higgs search channels at the Tevatron. Essentially, all analysis tools developed for the single top search can be used for this Higgs search.

2 Event Selection

Our single top event selection exploits the kinematic features of the signal final state, which contains a real $W$ boson, one or two bottom quarks, and possibly additional jets. To reduce multi-jet backgrounds, the $W$ originating from the top quark decay is required to have decayed leptonically. We demand therefore a high-energy electron or muon ($E_T(e) > 20$ GeV, or $P_T(\mu) > 20$ GeV/c) and large missing transverse energy (MET) from the undetected neutrino $\text{MET} > 25$ GeV. Electrons are measured in the central and in the forward calorimeter, $|\eta| < 2.0$. Exactly two jets with $E_T > 15$ GeV and $|\eta| < 2.8$ are required to be present in the event. A large fraction of the backgrounds is removed by demanding at least one of these two jets to be tagged as a $b$-quark jet by using displaced vertex information from the silicon vertex detector. The secondary vertex tagging algorithm identifies tracks associated with the jet originating from a vertex displaced from the primary vertex indicative of decay particles from relatively long lived $B$ mesons. The backgrounds surviving these selections are $t\bar{t}$, $W$ + heavy-flavor jets, i.e. $W + b\bar{b}$, $W + c\bar{c}$, $W + c$ and diboson events $WW$, $WZ$, and $ZZ$. Instrumental backgrounds originate from mis-tagged $W$ + jets events ($W$ events with light-flavor jets, i.e. with $u$, $d$, $s$-quark and gluon content, misidentified as heavy-flavor jets) and from non-$W$ + jets events (multi-jet events where one jet is erroneously identified as a lepton).

3 Background Estimate

Estimating the background contribution after applying the event selection to the single top candidate sample is an elaborate process. NLO cross section calculations exist for diboson and $t\bar{t}$ production, thereby making the estimation of their contribution a relatively straightforward process. The main background contributions are from $W + b\bar{b}$, $W + c\bar{c}$ and $W + c +$ jets, as well as mis-tagged $W$ + light quark jets. We determine the $W$ + jets normalization from the data and estimate the fraction of the candidate events with heavy-flavor jets using ALPGEN Monte Carlo samples, which were calibrated against multi-jet data. The probability that a $W$ + light-flavor jet is mis-tagged is parameterized using large statistics generic multi-jet data. The instrumental background contribution from non-$W$ events is estimated using side-band data with low missing transverse energy, devoid of any signal, and we subsequently extrapolate the contribution into the signal region with large missing transverse energy. The expected signal and background yield in the $W + 2$ jet sample is shown in Table 1 and graphically as a function of $W$ + jet multiplicity next to the table. Our analysis is performed in the $W + 2$ jet candidate sample with at least one tagged $b$-jet. This data sample features the highest signal purity. Table 1, however, demonstrates that the expected amount of single top events is much less compared to the large amount of expected backgrounds. In fact, the uncertainty on the backgrounds is larger than the expected signal, which renders a simple counting experiment impossible. This
Table 1: (Left) Expected signal and background yield in the $W + 2$ jet sample where at least one jet is tagged as a $b$-jet. (Right) Graphical sample composition as a function of the $W +$ jets multiplicity.

is exactly the reason why the search for single top quark production requires the best possible discrimination between signal and background processes and motivates the use of multivariate analysis tools.

3.1 Neural Network Jet-Flavor Separation

mistags and $W +$ charm events are a large class of background where no real $b$-quark is present and amount to about 50% of the $W + 2$ jets data sample even after imposing the requirement that one jet is identified by the secondary vertex $b$-tagger. We use a neural network tool which uses secondary vertex tracking information to distinguish $b$-jet events from charm and light-flavor jet events. Figure 2 shows the distribution of this jet-flavor separating neural network for the 644 $W + 2$ jets candidate events. All three single top analyses use this neural network tool to improve the sensitivity of the analyses.

4 Analysis Techniques

No single kinematic distribution encodes all conceivable signal background separation. We use several sophisticated analysis techniques to combine information into a single discriminant distribution which is used to extract the single top content in data.

4.1 Signal Significance and Discovery Potential

To quantify the significance of a potential single top signal we use the CLs/CLb method developed at LEP\textsuperscript{7}. In this approach, pseudo-experiments are generated from background only events (without single top) and from signal plus background events. We calculate the probability (p-value) of the background only pseudo-experiments to fluctuate to the observed result in data. A-priori we quote the expected sensitivity to a single top signal as the median p-value obtained from the signal + background pseudo-experiments. I.e. the quantity represents 50% of the generated signal plus background pseudo-experiments that had a p-value equal or greater than the expected p-value. All sources of systematic uncertainty are included in our statistical treatment and we consider correlation between normalization and discriminant shape changes due to sources of systematic uncertainty (e.g. the jet-energy-scale uncertainty).
4.2 Multivariate Likelihood Function Analysis

The multivariate likelihood function analysis computes a joint probability that a given candidate event originates from signal or background processes given a set of event characteristics $x_1, ..., x_{nvar}$. The likelihood ratio, as given in Equation 1, is used to build a likelihood function discriminant for s-channel and t-channel single top.

$$L(x_1, ..., x_{nvar}) = \frac{\prod_{i=1}^{nvar} p_{sig}^i}{\prod_{i=1}^{nvar} p_{bkg}^i} \quad \quad p_{sig}^i = \frac{N_{sig}^i}{N_{sig}^i + N_{bkg}^i}$$

The t-channel likelihood function discriminant is shown in Figure 2 which used seven input variables including the jet-flavor separating neural network described in the previous sub-section, the mass of the lepton, neutrino and the tagged jet $M_{\ell\nu b}$, a Matrix Element calculation from MadEvent and other event kinematic variables. The corresponding s-channel discriminant uses six input variables. The best expected p-value of 2.3% (2.1σ) is achieved by combining the t-channel and s-channel to a combined single top signal. The observed data show no indication of a single top signal and are compatible with a background-only hypothesis (p-value 58.5%). The upper limit on the combined single top cross section is 2.7 pb at the 95% C.L. The best fit for the combined s-channel and t-channel cross sections yields $\sigma_{s+t} = 0.3^{+1.2}_{-0.3}$ pb.

4.3 Neural Network Analysis

The neural network analysis combines 26 kinematic or event shape related variables to a discriminant output between -1 for background-like events, and +1 for signal-like events. The five most important input variables are reported by the neural network package. They are: the output of the jet-flavor separator neural network, $M_{\ell\nu b}$, the dijet mass, the pseudo-rapidity of the untagged jet multiplied by the charge of the detected lepton and the multiplicity of soft jets in the event with $8 \text{ GeV} < E_T < 15 \text{ GeV}$. To separate $t$- from s-channel single top quark production, two additional networks are trained and a simultaneous fit to both discriminants is performed. The combined search features an expected p-value of 0.56% (2.6σ). The observed p-value is 54.6%, providing no evidence for single top production. The corresponding upper limit on the cross section is 2.6 pb at the 95% C.L. The best fit yields $\sigma_{s+t} = 0.9^{+1.2}_{-0.5}$ pb.

4.4 Matrix Element Analysis

Using the Matrix Element analysis technique, we compute event-by-event probability densities that a given candidate event resulted from a given underlying interaction (signal or background
hypothesis). The measured four-vectors of the observed jets and the charged lepton serve as experimental input. The probability density is computed by integrating over the parton-level differential cross section \( d\sigma \), which includes the leading order matrix element for the process (calculated using MadEvent\(^8\)), the parton distribution functions \( f(x_i) \), and the detector resolutions parameterized by transfer functions \( W(y,x) \). Lepton momenta and jet angles are assumed well measured while the jet energy measurements are corrected to parton level energies using jet-energy to parton-energy transfer functions. We integrate over the quark energies and over the z-momentum of the neutrino to create a final probability density.

\[
P(x) = \frac{1}{\sigma} \int d\sigma(y) dq_1 dq_2 f(x_1) f(x_2) W(y,x)
\]

We use these probability densities to construct a discriminant variable for each event (Equation 3). We also introduce extra non-kinematic information by using the output \( b \) of the neural network jet-flavor separator which assigns a probability \( (0 < b < 1) \) for each \( b \)-tagged jet of originating from a \( b \) quark.

\[
EPD = \frac{b \cdot P_{\text{single top}}}{b \cdot P_{\text{single top}} + b \cdot P_{Wt} + (1-b) \cdot P_{WZ} + (1-b) \cdot P_{WZJ}}
\]

The expected p-value of the combined search is 0.6% (2.5\( \sigma \)). The observed p-value is 1.0% (2.3\( \sigma \)), providing a hint for a single top signal. The best fit yields \( \sigma_{s+} = 2.7^{+1.3}_{-1.8} \) pb.

5 Search for Heavy \( W' \) Resonances

Using the same dataset, we have searched for non-Standard Model production of single top quarks through a heavy \( W' \) boson resonance, \( pp \rightarrow W' \rightarrow t\bar{b} \rightarrow Wjj \) that appear in models with left-right symmetry, extra dimensions, Little Higgs, and topcolor\(^9\). We look for unexpected structure in the spectrum of the invariant mass of the reconstructed \( W \) boson and two leading jets \( (m_{Wjj}) \). No evidence for resonant \( W' \) production is observed and we exclude at the 95% C.L. a \( W' \) with Standard Model coupling strength and masses of 760 GeV/c\(^2\) (790 GeV/c\(^2\)) in case the right handed neutrino is smaller (larger) than the mass of the \( W' \) boson. These new limits exceed similar searches performed by CDF in Run I and D0 in Run II of the Tevatron program\(^10\).
6 Conclusions

We have performed searches for electroweak single top quark production at the Tevatron using 955 pb$^{-1}$ of data collected with the CDF II detector. The sensitivities to a single top signal at the rate predicted by the Standard Model range between 2.1 $\sigma$ and 2.6 $\sigma$ and are the most sensitive to-date. The multivariate likelihood function and neural network analysis observe a deficit of single top-like events in the data. The matrix element analysis observes a 2.3 $\sigma$ single top excess consistent with the Standard Model expectation. Using pseudo-experiment techniques, we estimated the compatibility of the three analyses to about 1.2 $\%$ given the correlation of about 60$\%$ and 70$\%$. Extensive cross-checks have been performed to understand the different outcomes in data. At present, there is no evidence for the cause other than statistical fluctuations given that the analyses work in different ways and make different, analysis specific assumptions. The larger datasets of 2000 pb$^{-1}$, already available to the CDF experiment, will clarify what the data are trying to tell us.

Using the same dataset, we exclude at the 95 $\%$ C.L. non-Standard Model single top quark production through heavy $W'$ resonances with masses of 760 GeV/c$^2$ (790 GeV/c$^2$) in case the right handed neutrino is smaller (larger) than the mass of the $W'$ boson.

Acknowledgments

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References

Evidence for single top quark production at DØ

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The results of the first analysis to show evidence for production of single top quarks are presented. Using 0.9 fb⁻¹ of data collected with the DØ detector at the Fermilab Tevatron, the analysis is performed in the electron+jets and muon+jets decay modes, taking special care in modeling the large backgrounds, applying a new powerful b-quark tagging algorithm and using three multivariate techniques to extract the small signal in the data. The combined measured production cross section is 4.8 ± 1.3 pb. The probability to measure a cross section at this value or higher in the absence of a signal is 0.027%, corresponding to a 3.5 standard deviation significance.

1 Introduction

The top quark is a very special fermion: it is by far the heaviest elementary particle discovered so far. Its large mass gives it a prominent role in the mechanism of electroweak symmetry breaking, since the Higgs boson coupling to fermions is proportional to their mass, and hence top quarks and Higgs bosons have a coupling strength of order unity. Top quarks also decay before they hadronize and thus pass their kinematic properties to the decay products, that can be then used to study the nature of “bare” quarks. The study of this elusive quark has provided us with some new insights into its properties recently, like its charge of 2/3e, its dominant decay to a W boson and a bottom quark, and its production cross section and mass, these last two being measured with ever increasing precision since its discovery.

Top quarks were first observed produced in t̄t pairs via the strong interaction at the Tevatron collider in 1995. In the SM, top quarks are also expected to be produced via the exchange of a W boson in s- or t-channel. The final state in these channels thus consists of one “single” top quark together with a b quark in the s-channel (t̄b) and an additional light quark in the t-channel (tq̄b). Single top quarks can also be produced in association with a W boson (tW̄), but the cross
section for this process at the Tevatron is very small and will be ignored here. The next-to-leading order prediction for the s-channel single top quark cross section is \( \sigma(p\bar{p} \to tb + X) = 0.88 \pm 0.11 \text{ pb} \), and for the t-channel process, the prediction is \( \sigma(p\bar{p} \to t\bar{q}b + X) = 1.98 \pm 0.25 \text{ pb} \).

Both the CDF and DØ collaborations have performed searches for this process in the past. In Run II, the best published 95% C.L. upper limits are \( \sigma(p\bar{p} \to tb + X) < 6.4 \text{ pb} \) and \( \sigma(p\bar{p} \to t\bar{q}b + X) < 5.0 \text{ pb} \). The analysis presented here, described in more detail in Ref., draws many techniques and experience from the previous DØ analyses, where it was made clear that multivariate techniques are necessary to be sensitive to the SM production cross section with limited data statistics.

2 Event selection

The data used in this search was collected from 2002 to 2005 with triggers that required an electron or a muon and at least one jet. The average trigger efficiency for e+jets events is 86% and 87% in the t- and s-channels, and for \( \mu + \text{jets} \), 82% and 87% in t- and s-channels signal samples respectively. Events are required to have exactly one isolated electron (muon) with \( p_T > 15 \text{ GeV} \) (18 GeV) within \( |\eta| < 1.1 \) (2.0), and \( E_T > 15 \text{ GeV} \). Events are also required to contain two, three or four jets, using a cone algorithm with radius \( R = \sqrt{(\Delta y)^2 + (\Delta\phi)^2} = 0.5 \) (where \( y \) is rapidity and \( \phi \) is azimuthal angle) to cluster energy deposits in the calorimeter. The leading jet has \( p_T > 25 \text{ GeV} \) and \( |\eta| < 2.5 \), the second leading jet has \( p_T > 20 \text{ GeV} \) and \( |\eta| < 3.4 \), and subsequent jets have \( p_T > 15 \text{ GeV} \) and \( |\eta| < 3.4 \).

The selection requirements described above achieve a summed signal to background ratio of around 1:180, with an acceptance of around 5.0% and 4.5% for s- and t-channel events respectively. To enhance the signal content of the selection, at least one jet in the event is required to be identified as a b-quark jet. A neural network has been used to identify jets originating from long-lived b hadrons. The variables used to identify such jets rely on the presence and characteristics of a secondary vertex and tracks with high impact parameters inside the jet. These variables are, ranked in order of separation power: (i) decay length significance of the Secondary Vertex Tagger (SVT); (ii) weighted combination of the tracks' IP significances; (iii) JIPIP probability (that the jet originates from the primary vertex based on the Jet Lifetime Probability algorithm); (iv) \( \Delta R \) per degree of freedom of the SVT secondary vertex; (v) number of tracks used to reconstruct the secondary vertex; (vi) mass of the secondary vertex; and (vii) number of secondary vertices found inside the jet. For a 0.5% light-jet (mis)tag rate, we obtain a 50% average tag rate in data for b jets with \( |\eta| < 2.4 \).

By requiring events to have at least one b-tagged jet, the signal to background ratio is enhanced to 1:22 on the sum of all channels, and the most sensitive channel, with two jets and one b-tagged jet, reaches a signal to background ratio of 1:10. The acceptance after b-tagging is reduced to 3.2% in the s-channel and 2.1% in the t-channel. The final event yields after b-tagging are given in Table 1, shown separated only by jet multiplicity. Figure 1 shows the data-background agreement in six basic distributions for all the channels combined.

3 Background normalization and agreement with data

Looking at Tab. 1, it is evident that if we are to reliably extract such a small signal from such large backgrounds (around 62 expected signal events from a total of 1,398 data events, where the error on the total background prediction is larger than the expected signals) we need to make sure the backgrounds are properly modeled in both their kinematics and overall normalization.

The \( t\bar{t} \) background is normalized to the integrated luminosity times the predicted \( t\bar{t} \) cross section of \( 6.8^{+0.8}_{-0.8} \text{ pb} \). The W+jets background, together with the multijet background,
Table 1: Numbers of expected and observed events in 0.9 fb⁻¹ for e and µ, 1 b tag and 2 b tag channels combined. The total background uncertainties are smaller than the component uncertainties added in quadrature because of anticorrelation between the W+jets and multijet backgrounds resulting from the background normalization procedure.

<table>
<thead>
<tr>
<th>Source</th>
<th>2 jets</th>
<th>3 jets</th>
<th>4 jets</th>
</tr>
</thead>
<tbody>
<tr>
<td>$ttb$</td>
<td>16 ± 3</td>
<td>8 ± 2</td>
<td>2 ± 1</td>
</tr>
<tr>
<td>$tqb$</td>
<td>20 ± 4</td>
<td>12 ± 3</td>
<td>4 ± 1</td>
</tr>
<tr>
<td>$tt\rightarrow\ell\ell$</td>
<td>39 ± 9</td>
<td>32 ± 7</td>
<td>11 ± 3</td>
</tr>
<tr>
<td>$t\bar{t}\rightarrow\ell\nu$+jets</td>
<td>20 ± 5</td>
<td>103 ± 25</td>
<td>143 ± 33</td>
</tr>
<tr>
<td>$Wbb$</td>
<td>261 ± 55</td>
<td>120 ± 24</td>
<td>35 ± 7</td>
</tr>
<tr>
<td>$Wc\bar{c}$</td>
<td>151 ± 31</td>
<td>85 ± 17</td>
<td>23 ± 5</td>
</tr>
<tr>
<td>$Wjj$</td>
<td>119 ± 25</td>
<td>43 ± 9</td>
<td>12 ± 2</td>
</tr>
<tr>
<td>Multijets</td>
<td>95 ± 19</td>
<td>77 ± 15</td>
<td>29 ± 6</td>
</tr>
<tr>
<td>Total background</td>
<td>686 ± 41</td>
<td>460 ± 39</td>
<td>253 ± 38</td>
</tr>
<tr>
<td>Data</td>
<td>697</td>
<td>455</td>
<td>246</td>
</tr>
</tbody>
</table>

Figure 1: Data-background agreement distributions for all channels combined for the lepton $p_T$, $E_T$, leading jet $p_T$, second leading jet $p_T$, the invariant mass of all jets, and the invariant mass of the reconstructed W together with the b-tagged jet. The hashed bands show the ±1 standard deviation uncertainty on the background.
normalized to the data in each channel (defined by lepton flavor and jet multiplicity) before $b$-tagging.

A second normalization to data is performed in order to determine how much of the total $W+\text{jets}$ sample is actually made of heavy-flavor jets, i.e. the fraction of $Wbb$ and $W\ell\ell$ in the total $W+\text{jets}$ sample. $W\ell\ell+\text{jets}$ is included in our model together with other $W+\text{light-quark jets}$, i.e.: inside $Wjj$ in Tab. 1. The heavy-flavor fraction is determined in our selected data on those events in which the neural network $b$-tagger fails to find a $b$ jet of the required quality. This sample of events is not used for the signal measurement and has approximately the same heavy-flavor composition as the signal sample. We find that a constant scale factor of $1.5\pm0.45$ applied to the $Wbb$ and $W\ell\ell$ components is necessary to achieve a good description of the data. The uncertainty assigned to this factor covers the expected dependence on event kinematics and the assumption that the scale factor is the same for $Wbb$ and $W\ell\ell$.

This constant heavy-flavor scale factor which absorbs higher order effects, was found to give overall a very good description of the data. It was checked that the most sensitive variables had well described shapes, specifically those distributions expected to suffer the largest shape dependence from higher order corrections, like the invariant mass of the two leading jets and the $p_T$ of the $b$-tagged jet. Two control samples, one enriched in $W+\text{jets}$ events and the other in $tt$ events, were also used to check for shape disagreements in the background model description of the data, and overall good agreement was found.

4 Systematic errors

The dominant contributions to the uncertainties on the backgrounds come from: normalization of the $tt$ background (18%), which includes a term to account for the top quark mass uncertainty; normalization of the $W+\text{jets}$ and multijet backgrounds to data (17–27%), which includes the uncertainty on the heavy-flavor fraction of the model; and the $b$-tagging probabilities (12–17% for double-tagged events). The uncertainty on the integrated luminosity is 6%; all other sources contribute at the few percent level. The uncertainties from the jet energy scale corrections and the $b$-tagging probabilities affect both the shape and normalization of the simulated distributions.

The 30% error assigned to the heavy-flavor fraction is by far the dominant uncertainty in the final cross section measurement. Its impact is nevertheless not directly 30% because it only affects two sources of the background (before $b$-tagging is applied) and it gets further reduced when taking into account the first normalization of $W+\text{jets}$ and QCD to data, such that in the end, the total uncertainty on the sum of $W+\text{jets}$ and QCD yields is between 17 and 27% (depending on the channel), including the heavy-flavor factor uncertainty.

5 Final separation methods and cross section measurements

Three separate multivariate techniques have been employed in DØ to extract the signal from the data: Boosted Decision Trees, Matrix Elements discriminants and Bayesian Neural Networks. The output distributions in the most sensitive channel, shown in Fig. 2 for the three methods, also demonstrate good overall agreement between data and background.

In order to extract the maximum information from the discriminant outputs, instead of cutting on the outputs and counting events, the full distributions are fed into a Bayesian statistical analysis to measure the single top quark production cross section. The expected and observed cross section results are summarized in Fig 3. The uncertainties include statistical and systematic components combined. The data statistics contribute 1.2 pb to the total 1.4 pb uncertainty on the $tb+tqg$ cross section for the DT analysis. The significance is measured re-running the analysis on 70,000 pseudo-datasets generated with all the uncertainties on the background model taken into account, but including only background sources. Thus we obtain the probability for
the background-only hypothesis to fluctuate up to give the measured (or SM) value of the \(t\bar{b}+tq\bar{b}\) cross section or greater.

To verify the kinematics of the selected events in the most discriminant region of the DT outputs, the \(t\)-channel characteristic distribution of \(Q^2(|\text{lepton}|) \times \eta(\text{untagged leading jet})\) is plotted in different slices of the DT output in Fig. 4. Requiring higher DT outputs clearly selects more signal-like data and the distinct asymmetric signal shape can be seen taking form.

The three analyses are highly correlated since they all use the same signal and background models and data, with almost the same systematic uncertainties. The correlation between the three methods has been measured in fake pseudo data-sets (which include the systematic uncertainties on our background model), with the SM single top cross section. The best linear unbiased estimator (BLUE) has been applied to the three measured values and their correlations to give a combined measured cross section of \(\sigma(pp \to t\bar{b} + X, \ tq\bar{b} + X) = 4.8 \pm 1.3 \text{ pb}\), which corresponds to a significance of 0.027\% (or 3.5 standard deviations).

![Figure 2](image1.png)

**Figure 2:** Output distributions in the electron channel with two jets, one of them \(b\)-tagged, for the Decision Tree discriminant (left), for the Matrix Element \(tq\) discriminant (center), and for the Bayesian Neural Net discriminant (right).

<table>
<thead>
<tr>
<th>(\sigma(tb\ + \ tq\overline{b}) [\text{pb}])</th>
<th>Expected</th>
<th>Observed</th>
<th>Significance</th>
</tr>
</thead>
<tbody>
<tr>
<td>DT</td>
<td>2.7^{+1.4}_{-1.4}</td>
<td>4.9 \pm 1.4</td>
<td>2.1\sigma</td>
</tr>
<tr>
<td>ME</td>
<td>3.0^{+1.5}_{-1.5}</td>
<td>4.6^{+1.8}_{-1.5}</td>
<td>1.8\sigma</td>
</tr>
<tr>
<td>BNN</td>
<td>3.2^{+2.0}_{-1.8}</td>
<td>5.0 \pm 1.9</td>
<td>1.3\sigma</td>
</tr>
</tbody>
</table>

![Figure 3](image2.png)

**Figure 3:** The observed results including the combination (left); and the expected and observed cross sections and significance for the three different multivariate analyses (right).

6 Measurement of \(|V_{tb}|\)

We use the decision tree measurement of the \(t\bar{b}+tq\bar{b}\) cross section to derive a first direct measurement of the strength of the \(V - A\) coupling \(|V_{td}f_t^L|\) in the \(Wtb\) vertex, where \(f_t^L\) is an arbitrary left-handed form factor. We measure \(|V_{td}f_t^L| = 1.3 \pm 0.2\). This measurement assumes \(|V_{td}|^2 + |V_{td}|^2 \ll |V_{td}|^2\) and a pure \(V - A\) and CP-conserving \(Wtb\) interaction. Assuming in addition that \(f_t^L = 1\), we obtain \(0.68 < |V_{tb}| \leq 1\) at 95% C.L.. These measurements make no assumptions about the number of quark families or CKM matrix unitarity.
7 Summary

To summarize, we have performed a search for single top quark production using 0.9 fb\(^{-1}\) of data collected by the D\(\bar{O}\) experiment at the Tevatron collider. We find an excess of events over the background prediction in the high discriminant output region from three analyses and interpret it as evidence for single top quark production. The excess has a combined significance of 3.5 standard deviations and the combined first measurement of the single top quark cross section is: \(\sigma(p\bar{p} \rightarrow tb + X, t\bar{q}b + X) = 4.8 \pm 1.3\) pb.

Acknowledgments

I thank the organizers of the Rencontres de Moriond for the enjoyable and fruitful atmosphere of the meeting; and the Marie Curie program of the European Union for funding my accommodation at La Thuile.

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Lepton distribution in top decay: A probe of new physics and top-polarization

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Abstract
We investigate the possibilities of studying new physics in various processes of $t$-quark production using kinematical distributions of the secondary lepton coming from decay of $t$ quarks. We show that the angular distributions of the secondary lepton are insensitive to the anomalous $tbW$ vertex and hence are pure probes of new physics in a generic process of $t$-quark production. The effects of $t$ polarization on the distributions of the decay lepton are demonstrated for top-pair production process at a $\gamma\gamma$-collider mediated by a heavy Higgs boson.

Keywords: Top, polarization, anomalous top coupling
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1 Introduction

As a result of the large mass of top quark and the consequent large coupling to the longitudinal component of the $W$ boson, it decays prior to hadronization. Hence, the spin of the decaying top quark leaves its imprint on the kinematical distributions of the decay products: the $W$ boson and the $b$ quark. Possible new physics of Electro-weak symmetry breaking (EWSB) may alter $t$ quark coupling with the $W$ boson leading to changed decay width and distributions. Further, new physics may also appear in its production process potentially affecting kinematical distributions and possibly polarization. The simultaneous presence of new physics both in production and decay processes of top quark may complicate the analysis and it may become difficult to probe new physics couplings of top quark. However, if one can construct observables that are sensitive to production and decay mechanism independent of each other, the analysis can be greatly simplified. It has been shown that the angular distribution of the leptons from the decay of top quark is insensitive to the anomalous $tbW$ coupling in the $2 \rightarrow 2$ processes of top quark production \cite{1} and also for a general process \cite{2}. Here we extend the result by relaxing the approximations used in the earlier analysis, for a general $2 \rightarrow n$ process\cite{3}.

2 Anomalous $tbW$ interaction and lepton distribution

We treat the $tbW$ vertex in a model independent way by parameterizing it as

$$\Gamma^\mu = \frac{-ig}{\sqrt{2}} \left[ \gamma^\mu (f_{1L}P_L + f_{1R}P_R) - \frac{i\sigma^{\mu\nu}}{m_W} (p_L - p_R)_{\nu} (f_{2L}P_L + f_{2R}P_R) \right]. \quad (1)$$
For the SM we have, \( f_{1L} = 1 \) and the anomalous couplings \( f_{1R} = f_{2L} = f_{2R} = 0 \). Further we note that the contribution from \( f_{1R} \) and \( f_{2L} \) is proportional to \( m_b \) and will be absent in vanishing \( m_b \) limit. To show the decoupling of lepton angular distribution from anomalous top coupling we assume,

1. \( t \)-quark is on-shell, i.e. using narrow-width approximation for top quark and

2. anomalous couplings \( f_{1R}, f_{2R} \), and \( f_{2L} \) are small, we retain terms linear in them.

Using narrow-width approximation for top quark we can write the matrix element square as

\[
|M|^2 = \frac{\pi \delta(p_t^2 - m_t^2)}{2 m_t \Gamma_t} \sum_{\lambda, \lambda'} \rho(\lambda, \lambda') \Gamma(\lambda, \lambda') \quad \rho(\lambda, \lambda') = M^*_\rho(\lambda) M^\dagger_\rho(\lambda') \quad \Gamma(\lambda, \lambda') = M^T(\lambda) M^T_\lambda(\lambda').
\]  

(2)

With this the differential cross section for \( 2 \rightarrow n \) process of top production followed by its decay can be written as,

\[
d\sigma = \sum_{\lambda, \lambda'} \left[ \frac{(2\pi)^4}{4} \Gamma(\lambda, \lambda') \delta^4(k_A + k_B - p_t - \sum_{i=1}^{n-1} p_i) \frac{d^3 p_t}{2E_t(2\pi)^3} \prod_{i=1}^{n-1} \frac{d^3 p_i}{2E_i(2\pi)^3} \right] 
\times \left[ \frac{1}{\Gamma_t} \left( \frac{(2\pi)^4}{2 m_t} \Gamma(\lambda, \lambda') \delta^4(p_t - p_b - p_\nu - p_\ell) \frac{d^3 p_b}{2E_b(2\pi)^3} \frac{d^3 p_\nu}{2E_\nu(2\pi)^3} \right) \frac{d^3 p_\ell}{2E_\ell(2\pi)^3} \right].
\]  

(3)

The first term in Eq. (3) can be written as \( d\sigma_{2-n}(\lambda, \lambda') dE_\ell d\cos \theta_\ell d\phi_\ell \) after integration in the lab frame, while the second term we integrate in the rest frame of the decaying top quark as:

\[
\frac{1}{\Gamma_t} \left( \frac{(2\pi)^4}{4} \frac{d^3 p_t}{2E_t(2\pi)^3} \frac{d^3 p_b}{2E_b(2\pi)^3} \frac{d^3 p_\nu}{2E_\nu(2\pi)^3} \Gamma(\lambda, \lambda') \delta^4(p_t - p_b - p_\nu - p_\ell) \right) = \frac{1}{32 \pi^2 m_t (2\pi)^4} \frac{E_\ell}{m_t E_\ell} \frac{d\Gamma}{dE_\ell d\cos \theta_\ell d\phi_\ell} .
\]  

(4)

Here the angular brackets indicate averaging over azimuthal angle between \( b \) quark and decay lepton. In the rest frame of top quark we have,

\[
\langle \Gamma(\lambda, \lambda') \rangle = g^4 m_t E_\ell^0 \left| \Delta_W(p_W^0) \right|^2 A(\lambda, \lambda') \times F(E_\ell^0),
\]  

(5)

where \( \Delta_W(p_W^0) = \frac{1}{p_{W}^0 - m_W^2 + i m_W \Gamma_W} \) is \( W \)-propagator, \( A(\pm, \mp) = (1 \pm \cos \theta_\ell) \), \( A(\pm, \mp) = \sin \theta_\ell e^{\pm \phi_\ell} \) is pure angular factor and \( F(E_\ell^0) \), given as

\[
F(E_\ell^0) = \left[ (m_t^2 - m_b^2 - 2p_\ell \cdot p_\ell) \left| f_{1L} \right|^2 + \text{Re}(f_{1L} f_{2L}^* m_t m_\ell \frac{p_{W}^0}{m_W m_p}) \right] - 2 \text{Re}(f_{1L} f_{2L}^*) m_b m_\ell \frac{p_{W}^0}{m_W m_p} - \text{Re}(f_{1L} f_{1R}^*) m_b m_\ell \frac{p_{W}^0}{m_W m_p},
\]  

(6)

is angle independent factor that depends upon the lepton energy. It should be noted here that all the anomalous couplings appear only in \( F(E_\ell^0) \). Combining production and decay part, we have

\[
d\sigma = \frac{1}{32 \pi^2 m_t (2\pi)^4} \sum_{\lambda, \lambda'} \frac{d\sigma_{2-n}(\lambda, \lambda') \times g^4 A^{c.m.}(\lambda, \lambda')}{E_\ell F(E_\ell) dE_\ell d\phi_\ell} .
\]  

(7)
Further, we have $\Gamma_t \propto \int E_t \, F(E_t) \, dE_t \, d\phi_t^{0W}$. Thus, if we integrate over $E_t$ in Eq. (7), the factor containing anomalous couplings cancels between the numerator and $\Gamma_t$ in the denominator. This proves that the differential rates are independent of anomalous $tbW$ couplings once we integrate over lepton energy. Hence, the angular distribution of lepton is independent of anomalous $tbW$ couplings in any inertial frame of reference and thus sensitive only to the production mechanism and possible new physics therein. On the other hand, the energy distribution of the decay lepton in the lab frame depends on the $tbW$ vertex, i.e. possible new physics in top decay, and also on energy-angular distribution of the produced top quarks, i.e. possible new physics in top production. However, the $E_t^{0W}$ distribution in the rest frame of the decaying top quark is given by $\frac{dN}{dE_t^{0W}} \propto \int E_t^{0W} F(E_t^{0W}) \, d\phi_t^{0W}$, hence sensitive only to the decay vertex with proportionality constant absorbing the information about production mechanism. In other words, the $E_t^{0W}$ distribution (in the rest frame of top quark) provides sensitivity to the anomalous $tbW$ coupling independent of the production mechanism of top quark. To summarize, the decay lepton distribution can provide separate and pure probes of possible new physics in production and decay of top quarks.

### 3 Polarization of top quark

Denoting the first line of Eq. (3) by $\sigma(\lambda, \lambda')$, one can write the polarization density matrix of the top-quark as

$$ P_t = \frac{1}{2} \begin{pmatrix} 1 + \eta_3 & \eta_1 - i \eta_2 \\ \eta_1 + i \eta_2 & 1 - \eta_3 \end{pmatrix}, $$

where

$$ \eta_3 = \frac{(\sigma(+,+) - \sigma(-,-))}{\sigma_{tot}}, \quad \eta_1 = \frac{(\sigma(+-) + \sigma(-,+))}{\sigma_{tot}}, \quad i \eta_2 = \frac{(\sigma(+,-) - \sigma(-,+))}{\sigma_{tot}} \quad (8) $$

and $\sigma_{tot} = \sigma(++,+) + \sigma(-,-)$. Here $\eta_3$ is the longitudinal polarization of top quark or the average helicity, $\eta_1$ is transverse polarization normal to the production plane and $\eta_1$ is transverse polarization in the production plane. These polarizations can be calculated from the angular distribution of leptons as:

$$ \frac{\eta_i}{2} = \frac{\sigma(p_L.s_i < 0) - \sigma(p_L.s_i > 0)}{\sigma(p_L.s_i < 0) + \sigma(p_L.s_i > 0)}, \quad \text{for } i = 1, 2, 3. \quad (9) $$

where $s_i$ and $p_L.s_i = 0$. For $p_t^u = E_t (1, \beta_t \sin \theta_t, 0, \beta_t \cos \theta_t)$, we have $s_1^L = (0, -\cos \theta_t, 0, \sin \theta_t)$, $s_2^L = (0, 0, 1, 0)$, $s_3^L = E_t (\beta_t \sin \theta_t, 0, \cos \theta_t)/m_t$. The measurement of $\eta_i$ requires knowledge of $s_L$. In other words, one requires to fully or partially reconstruct top momentum in the lab frame. Alternatively, one can look at the azimuthal distribution of leptons w.r.t. the top-production-plane as (at least) a qualitative probe of top polarization. For demonstration, we choose $\gamma\gamma \rightarrow tt$ with and without Higgs mediation and polarized photon beams. With help of variation in the photon polarization and presence and absence of Higgs exchange we generate ensemble of top quarks with varying polarization and show the decay lepton azimuthal distribution in Fig. 1. We note that for this example we have $\eta_1 = \eta_2 = 0$ and the values of $\eta_3$ are indicated on the plot. We see that the distribution is peaked at the top production plane ($\phi_t = 0$, $2\pi$) for positively polarized top and the height of peak decrease as polarization changes from positive to negative value. For large negative value of polarization ($-0.83$) there is a dip in the distribution in place of the peak near the top production plane. This qualitative feature of the azimuthal distribution can be converted into quantitative measure after establishing the correlation between $\eta_i$ and height of the peak for a given process.
Figure 1: The distribution in the azimuthal angle of the lepton in the lab frame for different top polarizations.

To conclude, we show that the lepton angular distribution is sensitive to top polarization and production mechanism independent of anomalous $tbW$ couplings and $E_T^j$ distribution is sensitive to anomalous $tbW$ coupling independent of top production mechanism. Further, we show that azimuthal distribution of decay lepton in the lab frame can be used to probe polarization of decaying top quark.

References


Tevatron: Top quark production and properties

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The measurements of top quark properties and the top pair production cross section at the Tevatron constitute important tests of the Standard Model predictions and could potentially be sensitive to new physics. I present in this report the latest results from the CDF and DØ collaborations in a Run II data sample of integrated luminosity up to 1.1 fb⁻¹.

1 Introduction

Since the top quark discovery in 1995 by the CDF¹ and DØ² collaborations, an extensive program is being conducted at the Tevatron to study and characterize its properties. Due to its large mass, the top quark is believed to play a special role in physics beyond the Standard Model (SM). Precision measurements have the potential to reveal effects from new physics. Previous Run I measurements, performed in 100 pb⁻¹ data samples, were limited by low statistics. With the Run II well underway, more than 2 fb⁻¹ of delivered integrated luminosity and analyzed datasets already ten times larger than those of Run I, the field of top physics has entered a new and exciting era.

The SM top quark at the Tevatron, a p̅p collider with √s = 1.96 TeV, is mostly produced in pairs with a theoretical NLO cross section prediction of 6.7 ± 0.8 pb³ for a top mass of 175 GeV/c². Since it decays to a W boson and a b quark basically 100% of the time, the final states are defined based on the decay channel of the W boson. Topologies in which both W’s decay to a lepton-neutrino pair are classified as “dilepton”, while “lepton+jets” events are those where only one W decays leptonically. Dilepton and lepton+jets represent 5% and 30% of the t̅t decays, respectively. The rest of the branching ratio corresponds to the case where either both W’s decay hadronically (45%) or when one or both decay into a tau lepton (20%). Top quark properties and cross section measurements presented here are performed in the lepton+jets or dilepton datasets with integrated luminosities up to 1.1 fb⁻¹.
2 Production Cross Section Measurements

The top quark pair production cross section measurement is an important test of QCD predictions and is key to other properties studies as it provides a well understood sample of candidate top quark events. This measurement is performed with a variety of techniques and in different topologies, which not only serves as a cross check, but is interesting in the search of new physics, as this would impact each final state differently. A brief description of the most recent results performed at the Tevatron is presented next, followed by a summary of the status of this measurement at the DØ and CDF collaborations.

2.1 Dilepton Channel

Due to the branching fraction (5% for $e/\mu$ analyses), this sample is small. However, it is a clean sample with contamination predominantly composed by Drell-Yan production and multi-jet events where a jet is misidentified as a lepton. While many of the studies are performed on very clean data sets obtained by means of tight lepton requirements, some use looser criteria to reduce statistical uncertainties. An example of the latter, is one of the most recent cross section measurements performed by the CDF collaboration. Using a 1.07 fb$^{-1}$ data sample, a cross section of $\sigma_{ee} = 9.0 \pm 1.3^{+0.5}_{-0.5}(\text{stat}) \pm 0.5^{+0.5}_{-0.5}(\text{syst}) \text{ pb}^4$ is obtained by selecting events with two or more jets, significant missing transverse energy and two opposite-sign leptons, where one is identified as an isolated track.

DØ's most recent and precise measurement in this channel, utilizes an integrated luminosity of approximately 1 fb$^{-1}$ and combines the cross section obtained in events with $ee$, $e\mu$ and $\mu\mu$ final states. It requires 2 oppositely charge leptons, 2 or more jets for the $ee$ and $\mu\mu$ channels, one or more for $e\mu$ and a set of optimized cuts per final state to discriminate against background. The cross section in each individual channel is extracted by minimizing a negative log-likelihood function based on the Poisson probability to observed a certain number of events given the luminosity, branching fraction, efficiency and the number of background events. The combined result is obtained by minimizing the sum of the negative log-likelihood function for each individual final state and is found to be $\sigma_{ee} = 6.8^{+1.2}_{-1.1}(\text{stat})^{+0.9}_{-0.8}(\text{syst}) \pm 0.4(\text{lumi})$ pb$^{-5}$.

2.2 Lepton + Jets Channel

Although the lepton+jets sample benefits from its larger branching fraction, there is significant background contamination from W boson production in association with jets and non-W QCD events (instrumental background). Signal to background discrimination can be improved using techniques to identify b-jets, a process known as "b-tagging". The latest and most precise DØ measurement, performed on a data set of about 900 pb$^{-1}$, uses a neural network that combines the outputs of several b-tagging algorithms rendering an efficiency of 54% and a misidentification rate below 1%. The event selection requires an identified lepton, large missing transverse energy and three or more jets. This sample is divided into 8 channels based on the lepton flavor (electron or muon), the jet multiplicity (three and four or more) and the number of tagged jets (one and two or more). The cross section is then calculated by performing a maximum likelihood fit to the observed number of events in the various channels. DØ finds a cross section of $\sigma_{ee} = 8.3^{+0.6}_{-0.5}(\text{stat})^{+0.9}_{-0.8}(\text{syst}) \pm 0.5(\text{lumi})$ pb$^{-5}$. 
2.3 Summary

Figure 1 shows the summary\(^a\) of the most recent top quark pair production cross section measurements from the Tevatron. Both the DØ and CDF collaborations perform the measurement in different topologies with various techniques. Results are in agreement with each other and with the theoretical prediction. The most precise measurements are no longer statistically dominated and have uncertainties of \(\sim 14\%\). Shown is also the CDF combination result (analyses with integrated luminosity up to 760 pb\(^{-1}\)) which has an uncertainty of 12% comparable with the uncertainty from the theoretical prediction.

![Graph showing top quark pair production cross section measurements from DØ and CDF.

Figure 1: Summary of most recent top quark production cross section measurements from DØ (left) and CDF (right). The band represent the NLO theory calculation for a 175 GeV/c\(^2\) top quark mass.]

3 Top quark Properties

The goal of measuring the top quark properties is to find out whether the top quark behaves as the SM predicts or if new physics is present in the \(t\bar{t}\) data sample.

3.1 Production Mechanism

The Standard Model predicts that, at the Tevatron, about 85% of the top pair events are produced through quark-antiquark annihilation while approximately 15% are produced through gluon fusion. CDF tests this prediction using the low \(p_T\) charged particle multiplicity as the discriminator between the two different production channels and performs the first measurement of the ratio \(\sigma(gg \rightarrow t\bar{t})/\sigma(pp \rightarrow t\bar{t})\) using 1fb\(^{-1}\) of data. Figure 2 (left plot) shows the correlation between the average number of gluons \((N_g)\) and the average number of low \(p_T\) charged particles \((N_{\text{track}})\) in different samples. While \(N_{\text{track}}\) is measured in high \(p_T\) inclusive lepton or inclusive jet data samples, \(N_g\) is obtained from simulation, where only gluons which are part of the hard scattering matrix element are used in the calculation. The higher the number of gluons in the sample, the larger the low \(p_T\) track multiplicity is. Using no-gluon and gluon-rich low \(p_T\) track multiplicity distributions, calibrated on data, CDF fits the signal sample of tagged \(W+4\) (or more) jets events for the gluon rich fraction and obtains \(f_g = 0.07 \pm 0.15\text{(stat)} \pm 0.07\text{(syst)}\). Figure 2 presents the two components (gluon rich and no-gluon) and fit result. Taking into

\(^a\)CDF Dilepton and lepton+jets results in the 1.2 fb\(^{-1}\) dataset became available after the Conference.
account the relative acceptance between gg and q̅q events and the background contributions, a ratio of \( \sigma(gg \rightarrow t\bar{t})/\sigma(p\bar{p} \rightarrow t\bar{t}) = 0.01 \pm 0.16(\text{stat}) \pm 0.07(\text{syst}) \) is found, consistent with the SM expectation.

Figure 2: Left: Correlation between the average low \( p_T \) track multiplicity and the average number of gluons \( <N_g> \) for different data samples. Right: Fit result for the tagged W+4 or more jet data sample. The two components of the fit (gluon rich and no-gluon) are also shown.

3.2 Anomalous Production

Non-SM top pair production mechanisms, like \( t\bar{t} \) production via an intermediate resonance, are proposed by different theories. For example, the topcolor-assisted Technicolor\(^8\) model predicts a heavy \( Z' \) boson which couples preferentially to the third generation of quarks and with cross sections expected to be visible at the Tevatron. These processes may manifest as narrow structures in the total pair mass, \( M_{t\bar{t}} \), superposed on the SM expected spectrum.

Both the DØ and CDF collaborations have performed model independent searches for a narrow width heavy resonance in the lepton+jets channel by studying the \( t\bar{t} \) invariant mass spectrum, reconstructed using a constrained kinematic fit. Results are found to be consistent with the SM expectation. Modeling the resonance as a \( Z' \)-like boson, 95% CL upper limits are obtained for different resonances mass hypotheses. For a leptophbic Topcolor model with a resonance width of 1.2% of its mass and in a data sample of approximately 1 fb\(^{-1}\), CDF excludes masses up to 725 GeV/c\(^2\)\(^9\). For a data sample of 370 pb\(^{-1}\), DØ excludes masses below 680 GeV/c\(^2\)\(^10\). Figure 3 shows CDF \( M_{t\bar{t}} \) spectrum and upper limits.

Figure 3: Left: Reconstructed invariant mass of the \( t\bar{t} \) system. Right: Expected and experimental upper limits on the \( Z' \rightarrow t\bar{t} \) production cross section together with leptophbic topcolor \( Z' \) cross section prediction.
3.3 Examining the tWb vertex: W helicity

In the SM, top quark decays into a W boson and a b quark with a branching ratio close to 100%. The V-A coupling of the tWb vertex only allows the W boson to have longitudinal or left-handed polarizations while the right-handed polarization is suppressed. The fraction of decays to longitudinal and left-handed W bosons is expected to be $F_0^{SM} \sim 0.7^{11}$ and $F_- \sim 0.3$, respectively. Thus, measuring the W helicity provides a direct test of the V-A nature of the tWb coupling and a probe of new physics beyond the SM.

Through the study of sensitive observables, like lepton transverse momentum, invariant mass of the lepton and b-jet or $\cos \theta^*$ (angle between the lepton momentum in the W boson rest frame and the W momentum in the top quark rest frame), DØ and CDF collaborations have measured $F_0$ and the right-handed $F_+$ fractions in both the lepton+jets and dilepton channels. One of the most recent measurements done by CDF on 955 pb$^{-1}$ of data, uses lepton+jets final state events and $\cos \theta^*$ as the discriminant variable. To build the $\cos \theta^*$ distribution in data, the kinematics of the $t\bar{t}$ candidates is fully reconstructed. Performing a binned likelihood fit, with a fixed $F_+ = 0$ value, the longitudinal fraction is determined to be $F_0 = 0.39 \pm 0.12^{+0.07}_{-0.08}$ (syst). When fixing $F_0$ to its SM value of 0.7, a fraction $F_+ = -0.03 \pm 0.06$ (stat) $^{+0.04}_{-0.10}$ (syst) is obtained. An upper limit of $F_+ < 0.1$ is set at the 95% CL. All DØ and CDF results for the W boson helicity are in agreement with the SM prediction within large statistical uncertainties.

3.4 Top Quark Charge

One of the fundamental quantities that characterize the top quark is its electric charge. Due to the ambiguity in the assignment of the b quark and W boson to which the top decays, it is possible to reconstruct an object of charge -4/3, instead of the 2/3 top quark charge that the SM predicts. In fact, such an hypothesis has been proposed in which this new particle would correspond to an exotic quark, part of a fourth generation of quarks and leptons. Thus, determining whether the top quark (t) decays into a $W^-$ and a b quark ($W^{-} b$) would indirectly indicate that the top charge is indeed 2/3. Both CDF and DØ have performed such a measurement. Presented here is the CDF result using a data sample of 955 pb$^{-1}$ in the dilepton channel and 695 pb$^{-1}$ in the lepton+jets one. Double tagged lepton+jets events are reconstructed by means of a kinematic fitter while a discriminant based on the invariant mass of the lepton and the b-jet is used in single tagged dilepton ones. While the charge of the W boson is defined by the charge of the lepton, the flavor of the b-jets is assigned using a Jet Charge Algorithm. This algorithm, which sums up the charge of the tracks assigned to the jet, each weighted by its momentum along the jet axis, is calibrated in $b\bar{b}$ data events. Each W-jet pair is labeled as being standard model like (SM-like) or exotic quark model like (XM-like) based on the product of their charges, negative for the SM and positive in the XM case. In order to obtain a confidence limit on either hypothesis, a profile likelihood method is used to build a likelihood curve as a function of the fraction of SM-like events. Defining an a-priori probability of 1% ($\alpha$) to incorrectly reject the SM when this model is true, the sensitivity, defined as the probability of rejecting the SM in the case that the XM model is true, is found to be 81%.

CDF observes 62 SM-like pairs and 48 XM-like ones, which corresponds to a p-value of 0.35 under the SM hypothesis. Since this p-value is larger than the $\alpha$ value, the result is found to be consistent with the SM and the XM hypothesis is excluded at 81% C.L. Figure 4 shows the distribution of the product of the W charge and the jet charge for the 110 observed pairs.
Figure 4: Product of the W charge and the associated jet charge for Data and MC Lepton+Jets pairs (left) and Dilepton ones (right). Negative values correspond to SM-like pairs ($W^+b$ or $W^-b$).

4 Conclusion

I have presented the latest top quark properties and pair production cross section measurements from the Tevatron in Run II with an integrated luminosity of up to 1.1fb$^{-1}$. Cross section results, which have reached a level of precision comparable with the theoretical one, are consistent with the NLO QCD predictions and in agreement with each other. Several properties measurements have been reported. So far no evidence of new physics has been found. With the Tevatron and the CDF and DØ experiments performing very well, an exciting landscape in top physics is foreseen for the coming years.

Acknowledgments

I thank the organizers for a stimulating conference and my CDF and DØ colleagues for their efforts to produce the results presented in this work.

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Beyond SM, Extradim

B Physics
Tevatron Results on the Discovery of $\Sigma_b^{(*)}$, $B_s$ Oscillations and the Measurement of $\Delta m_s$, the Lifetime Difference $\Delta \Gamma_s$ and the CP-violating phase $\phi_s$.

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I discuss results from the Tevatron experiments on mixing and CP-violation in $B_s$ mesons, including the observation of $B_s$ oscillations and the first precision measurement of the mixing frequency, as well as a measurement of the lifetime difference $\Delta \Gamma_s$ and the first measurement of the CP-violating phase $\phi_s$. I also briefly report on the observation of four new bottom baryons at CDF.

1 Introduction

The Tevatron $pp$ collider at Fermilab, operating at a center of mass energy of $\sqrt{s} = 1.96$ TeV, is a unique place to study the properties of $B$ baryons and the heavy $B$ mesons: the $B_s$, $B_c$ and their excited states. By now, the two experiments, D0 and CDF have collected an integrated luminosity of $\sim 2.5$ fb$^{-1}$ each, of which the analyses discussed here use between 1.0 and 1.3 fb$^{-1}$.

The hadronic environment makes triggering an important prerequisite for $B$ physics. D0 collects events containing $B$ mesons by triggering on muons produced in semi-leptonic decays. The excellent muon coverage allows the collection of large samples of semileptonically decaying $B$ mesons, which drive the D0 analyses discussed here. In addition to lepton triggers, CDF exploits triggers which find tracks that are displaced from the primary $pp$ interaction point. These triggers give unique access to hadronic $B$-decays, which are of prime importance in both the CDF $B_s$ oscillation and the $\Sigma_b^{(*)}$ analyses.

2 Observation of $\Sigma_b^{(*)}$ at CDF

CDF has observed four new bottom baryons, $\Sigma_b^+$ and $\Sigma_b^{*-}$, decaying into $\Lambda_b^0\pi^\pm$, where $\Lambda_b^0 \to \Lambda_c^+\pi^-$ and $\Lambda_c^+ \to p^+K^0\pi^+$ (charge conjugation is implied). So far, only the lightest b baryon, the $\Lambda_b$ ($ud\bar{b}$) had been observed. The $\Sigma_b^{*+(-)}$ has quark content $buu$ ($b\bar{d}d$) and $J^P = \frac{1}{2}^+$, whereas
Figure 1: Distribution of the mass difference $Q$ for $\Lambda_b^0\pi^-$ (top) and $\Lambda_b^0\pi^+$ (bottom) combinations. The unbinned maximum likelihood fit and the fitted contributions of $\Sigma_b$ and $\Sigma_b^*$ are also shown.

The excited states $\Sigma_b^{*\pm}$ have $J^P = \frac{3}{2}$. Observation of these states, and the measurement of their mass splittings, provides a test of QCD, which is predictive in this case because of the large mass of the $b$ quark, through e.g. heavy quark effective theory. Figure ?? shows the distribution of the mass difference $Q = m(\Lambda_b^0\pi^-) - m(\Lambda_b^0) - m_\pi$. The four new states are clearly visible as excesses over the background. The significance of this signal has been determined to exceed 5.2$\sigma$. A more detailed description of this analysis can be found in J. Pursley's contribution to these proceedings and in $^7$.

3 Lifetime difference and CP-violation in the $B_s$ system.

The phenomenology of mixing between $B_s$ and $\bar{B}_s$ mesons is characterized by the mass difference between the mass eigenstates of the system $\Delta m_s$, which is equal to the $B_s - \bar{B}_s$ oscillation frequency, the difference in lifetime $\Delta\Gamma_s$ of the mass eigenstates, and the $CP$-violating phase $\phi_s$. The phase $\phi_s$ is predicted to be small in the Standard Model (SM) but new physics could introduce (large) $CP$-violation in the $B_s$ system. Observation of a sizable value of $\phi_s$ would therefore be an definite sign of physics beyond the standard model. The lifetime difference $\Delta\Gamma_s$ is related to $\phi_s$: $\Delta\Gamma_s = \Delta\Gamma_s^{SM} \times [\cos(\phi_s^{SM} + \phi_s^{NP})]$. The SM prediction is $\Delta\Gamma_s^{SM} = 0.096 \pm 0.039$ $\tau$, but a large $CP$ violating phase would lower this value, making the measurement of $\Delta\Gamma_s$ another important probe for new physics.

3.1 Charge Asymmetry measurements

$CP$-violation in $B_s$ mixing could give rise to a charge asymmetry in (semileptonically) decaying $B_s$ mesons:

$$A^s_{SL} = \tan(\phi_s) \frac{\Delta\Gamma_s}{\Delta m_s}. \quad (1)$$
D0 obtains a first handle on the on the change asymmetry in $B_s$ decays from a measurement of the same-sign dimuon charge asymmetry:

$$A_{SL}^{\mu\mu} = \frac{N(bb \rightarrow \mu^+\mu^+) - N(bb \rightarrow \mu^-\mu^-)}{N(bb \rightarrow \mu^+\mu^+) + N(bb \rightarrow \mu^-\mu^-)}.$$  

(2)

Here, one of the muon tags the initial flavor of the $B$ meson, while the other $\mu$ reveals the flavor at decay. A net positive (negative) asymmetry would imply that the $B^0 \rightarrow \bar{B}^0$ transition (yielding $\mu^+\mu^+$) occurs more (less) often than the charge conjugate process; a clear violation of $CP$. Using a sample of $\sim 310,000$ dimuon pairs, D0 obtains: $A_{SL}^{\mu\mu} = -0.0092 \pm 0.0044$(stat) $\pm 0.0032$(syst)\(^7\). This asymmetry has contributions from both neutral $B$ mesons: $B_d$ and $B_s$. Using world average values for $B_d$ and $B_s$ production fractions and the $B_s^0$ charge asymmetry measurements from the B-factories, the charge asymmetry in $B_s$ mixing is extracted\(^7\) as:

$$A_{SL}^{\mu\mu,B_s} = -0.0064 \pm 0.0101$(stat + syst).

(3)

In another measurement\(^7\), D0 have directly measured the untagged charge asymmetry in a sample of reconstructed $B_s$ decays into $\mu D_s$, with $D_s \rightarrow \phi \pi$ and $\phi \rightarrow K^+K^-$. Assuming equal cross-sections for the production of $B_s$ and $\bar{B}_s$, they find:

$$A_{SL}^{\mu D_s} = \frac{N(\mu^+D_s^0) - N(\mu^-D_s^+)\nu}{N(\mu^+D_s^0) + N(\mu^-D_s^+)} = 0.0245 \pm 0.00193$(stat) $\pm 0.0035$(syst).

(4)

In both asymmetry measurements, certain sources of systematic uncertainty are reduced by regular flipping of the magnetic fields in D0. The remaining systematic uncertainties in the two measurements are dominated by different effects and the overlap between the two samples is only $\sim 10\%$. The combination of the two results is straightforward and yields the best estimate of the charge asymmetry in semileptonic $B_s$ decays:

$$A_{SL}^{B_s} = -0.0001 \pm 0.0090$(stat + syst),

(5)

which is consistent with zero and with the SM prediction.

3.2 Lifetime difference of $CP$-violating phase.

The decay $B_s \rightarrow J/\psi \phi$, proceeding through the $b \rightarrow c\bar{c}s$ transition, gives rise to both $CP$-even and $CP$-odd final states, which can be separated in a simultaneous analysis of the time evolution and the distributions of three decay angles that convey the $CP$ information.

D0 have performed this measurement using a sample containing $\sim 23,000$ untagged signal events\(^7\). Assuming no large $CP$-violation, as predicted by the SM, the mass eigenstates and the $CP$ eigenstates coincide. Under this assumption, the fit yields:

$$\Delta \Gamma_s = \Delta \Gamma_{CP} = 0.12^{+0.08}_{-0.10} \pm 0.02$(syst) ps\(^{-1}\),

(6)

in accordance with the SM prediction, with an older measurement from CDF\(^7\), and with the value of $\Delta \Gamma_s$ obtained from the branching ratio $B_s \rightarrow D_s^{(*)}D_s^{(*)}\$\(^7\). In a more general fit to same data, $\phi_s$ is treated as a free parameter, allowing for possible $CP$-violation through contributions from new physics. In this case, the result is (up to a 4-fold ambiguity):

$$\Delta \Gamma_s = 0.17 \pm 0.09$(stat) $\pm 0.02$(syst) ps\(^{-1}\),

$$\phi_s = -0.79 \pm 0.58$(stat) $^{+0.14}_{-0.09}$(syst),

(7)

(8)

which constitutes the first direct measurement of the mixing phase.
Figure 2: The $\phi_s$ vs $\Delta \Gamma_s$ plane, showing the four degenerate best fit values for the unconstrained fit to the $B_s \to J/\psi \phi$ data and the $\Delta \ln(L)$ error ellipse (blue, dashed) for one of the SM-like solutions. The combined semileptonic charge asymmetry band (yellow) shows the constraint from Eq. (??). The red star and error contour indicate the result with the constraint from the semileptonic charge asymmetry measurements and from the world average flavor specific lifetime. A band representing the relation $\Delta \Gamma_s = \Delta \Gamma_s^{SM} \times |\cos \phi_s|$ for SM allowed range of $\Delta \Gamma_s$ is shown in green.

3.3 Combination

Since $\Delta m_s$ has been measured precisely (see next section), Eq. (??) can be used to constrain the measurement of $\Delta \Gamma_s$ and $\phi_s$ by using the charge asymmetry measurements.

Further constraints are provided by measurements of the $B_s$ lifetime using flavor-specific decays, which are measured by fitting a single lifetime, $1/\Gamma_{BS}$, to the decays. In case there are two components to the lifetime distribution (non-zero $\Delta \Gamma_s$), the fitted value of the flavor specific lifetime deviates from the average lifetime, $1/\bar{\Gamma}_s = \bar{\Gamma}_s - (\Delta \Gamma_s)^2/2\bar{\Gamma}_s + O((\Delta \Gamma_s)^3/\bar{\Gamma}_s^2)$. The world average flavor specific lifetime is thus used to further constrain the measurement of $\Delta \Gamma_s$ and $\phi_s$.

Figure ?? shows the constrained and unconstrained results. The combined measurement yields $\Delta \Gamma_s = 0.19 \pm 0.09$ ps$^{-1}$, which is in good agreement with the SM prediction, and $\phi_s = -0.70^{+0.47}_{-0.39}$, which is consistent at the 1.8$\sigma$ level with the SM prediction, $|\phi_{SM}| = (4.2\pm1.4) \times 10^{-3}$.

4 $B_s$ oscillations

Resolving the rapid oscillations between $B_s$ and $\bar{B}_s$ and making a precise measurement of the oscillation frequency $\Delta m_s$ have been important goals of the experiments at the Tevatron and elsewhere.

At Moriond 2006, D0 presented the first double sized bound on $\Delta m_s$, based on a sample of 27,000 $B_s \to \mu \nu D_s$ decays, with $D_s \to \phi \tau$. Since then, the analyses has been improved by including $D_s \to K^*K$ and $D_s \to K^*K$ and $B_s \to e\nu e D_s$, with $D_s \to \phi \tau$, bringing the total signal yield to 42,000 semileptonically decaying $B_s$ events.

CDF have used a sample of $\sim 62,000$ $B_s \to \mu \nu D_s$ and $e\nu e D_s$ decays, collected using both leptonic and displaced track triggers. The latter also give access to a sample of 5600 fully reconstructed hadronic decays, $B_s \to D_s \pi$ and $B_s \to D_s \pi \pi$, with $D_s \to \phi \tau$, $K^*K$ or $\pi \pi$. Another 3100 signal events are obtained in the channel $B_s \to D_s^\ast \pi$ and $B_s \to D_s \rho$, with
\( D_s \to \phi \pi \). In these modes a low momentum \( \pi^0 \) or \( \gamma \) escapes detection.

### 4.1 Lifetime resolution

In order to resolve the rapid oscillations, an accurate determination of the decay time of the \( B_s \) mesons is crucial. The experiments use their silicon trackers to accurately determine the flight distance in the transverse plain \( L_T \). The decay time is \( t = L_T \times m_b / e P_T \), where \( m_b \) is the \( B_s \) mass and \( P_T \) is the transverse momentum of the \( B_s \).

The semileptonic events, have an intrinsic \( P_T \) uncertainty due to the unobserved momentum of the neutrino. This translates in an uncertainty on the decay time, which is detrimental to the sensitivity for rapid oscillations (large \( \Delta m_s \)). For this reason, CDF's hadronic events are particularly valuable. Despite the relatively low number of events, they dominate in the signal. The average decay time resolution is about 87 (150) fs for CDF's hadronic (semileptonic) events and 160 fs in the D0 analysis. This should be compared to the \( B_s \) oscillation period of about 300 fs (\( \Delta m_s = 18 \) ps\(^{-1} \)).

### 4.2 Initial flavor tagging

While the flavor at decay is unambiguously identified from the decay products, the oscillation measurement also requires knowledge of the flavor at production. Opposite Side Taggers (OST) infer the \( B_s \) production flavor from the decay products of the other \( B \) hadron that results from the initially produced \( b \bar{b} \) pair. Various OSTs measure the charge of leptons from semileptonically decaying \( B \) hadrons, of kaons resulting from the \( b \to c \to s \) decay chain, or of the jet produced by the \( B \) hadron decay. CDF combines these OSTs using a neural network, while D0 uses a combination of jet charge and lepton tags.

In addition, CDF has developed a Same Side Kaon Tagger (SSKT). The idea behind this tagger is that an \( s \bar{s} \) quark pair must be created in the hadronization of a \( b \) quark into a \( B_s \). The remaining \( s \) quark will combine with a light quark into a kaon. The charge of the kaon identifies the production flavor of the \( B_s \). CDF uses its Time-of-flight detector, a measurement of \( dE/dx \) in the central tracker, and kinematic variables to identify kaons in a cone around the \( B_s \).

The performance of the flavor taggers is quantified by the efficiency \( \epsilon \) and the dilution \( D_D \), which is related to the wrong-tag probability \( w \), \( D_D \equiv 1 - 2w \). The effectiveness, \( \epsilon D_D^2 \) of the OSTs used in D0 and CDF is 2.5% and 1.8% respectively. CDF's SSKT has \( \epsilon D_D^2 = 3.5\% \) (hadronic) and 4.8% (semileptonic) and thus contributes most to the sensitivity of the CDF analysis.

### 4.3 Results

The most recent D0 amplitude scan is shown in Figure ??(left). The amplitude is consistent with unity around \( \Delta m_s = 18 \) ps\(^{-1} \). A 90% C.L. double sided confidence bound is obtained: \( 17 < \Delta m_s < 21 \) ps\(^{-1} \). The probability for a random fluctuation to produce such a signal is about 8%, as estimated from Monte Carlo.

After having published a first result \(^1\), CDF have upgraded their 1 fb\(^{-1} \) analysis \(^2\). The main improvements are the application neural networks for the signal selection and in the flavor taggers, and the inclusion of the partially reconstructed hadronic decays. The amplitude scan is shown in Figure ??(right). At its highest point, the amplitude is consistent with one and clearly incompatible with zero. The significance has been estimated from the data by randomizing the tagger decisions to create pseudo-experiments devoid of an oscillation signal. 28 out of \( 3.5 \times 10^6 \) trials were more signal-like than the main result, giving the signal a significance of 5.4 \( \sigma \). This signal allows for a precise measurement of \( \Delta m_s \):

\[
\Delta m_s = 17.77 \pm 0.10 \text{(stat)} \pm 0.07 \text{(syst)} \text{ ps}^{-1},
\]  
(9)
which is in good agreement with the SM expectation. This result has been used to determine the ratio of CKM matrix elements

$$\frac{|V_{td}|}{|V_{ts}|} = 0.2060 \pm 0.0007 (\text{stat})^{+0.0008}_{-0.0006} (\text{theo} + \Delta m_d). \quad (10)$$

5 Conclusions

The Tevatron experiments are combing the $B_s$ system for new physics. Both have seen $B_s \bar{B}_s$ oscillations and CDF has measured the mixing frequency with $\sim 1\%$ accuracy, using a signal with $> 5\sigma$ significance. The experiments have started to look for $CP$-violation in $B_s$ mixing. D0 has made charge asymmetry measurements and performed the angular analysis of $B_s \rightarrow J/\psi \phi$ decays, resulting in the first precision measurement of $\Delta \Gamma_s$ and the first determination of the $CP$-violating phase $\phi_s$. While based on only $\sim 50\%$ of the data currently on tape, these measurements already put powerful constraints on models of new physics (see the contribution of P. Ball to this conference). Significant improvements will be made as the Tevatron keeps producing data. In particular, the taggers developed for the $B_s$ oscillation analysis will be used to perform a tagged measurement in $B_s \rightarrow J/\psi \phi$, improving the sensitivity to $\Delta \Gamma_s$ and $\phi$.

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PROBING NEW PHYSICS THROUGH $B_s$ MIXING

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I discuss the interpretation of the recent experimental data on $B_s$ mixing in terms of model-independent new-physics parameters.

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1 Introduction

One of the most promising ways to detect the effects of new physics (NP) on $B$ decays is to look for deviations of flavour-changing neutral-current (FCNC) processes from their Standard Model (SM) predictions; FCNC processes only occur at the loop-level in the SM and hence are particularly sensitive to NP virtual particles and interactions. A prominent example that has received extensive experimental and theoretical attention is $B^0_q - \bar{B}^0_q$ mixing ($q \in \{d, s\}$), which, in the SM, is due to box diagrams with $W$-boson and up-type quark exchange. In the language of effective field theory, these diagrams induce an effective local Hamiltonian, which causes $B^0_q$ and $\bar{B}^0_q$ mesons to mix and generates a $\Delta B = 2$ transition:

$$\langle B^0_q | H^{\Delta B=2} | \bar{B}^0_q \rangle = 2M_{B_q} M_{12}^q,$$

where $M_{B_q}$ is the $B_q$-meson mass. Thanks to $B^0_q - \bar{B}^0_q$ mixing, an initially present $B^0_q$ state evolves into a time-dependent linear combination of $B^0_q$ and $\bar{B}^0_q$ flavour states. The oscillation frequency of this phenomenon is characterized by the mass difference of the "heavy" and "light" mass eigenstates,

$$\Delta M_q \equiv M^q_H - M^q_L = 2|M^q_{12}|,$$

and the CP-violating mixing phase

$$\phi_q = \arg M^q_{12},$$

which enters mixing-induced CP violation. While the mass difference in the $B_d$ system has been known for a long time, $\Delta M_s$ has only been measured in 2006, by the CDF collaboration, with the result $^1$

$$\Delta M_s = [17.77 \pm 0.10 \text{(stat)} \pm 0.07 \text{(syst)}] \text{ps}^{-1}.$$  

In Ref. $^2$, we have discussed the impact of this result on a model-independent parametrisation of NP in the $B_q$ system. In the meantime, experimental information has become available also for the mixing phase in the $B_s$ system $^{3,4,5,6}$. In these proceedings, we update the constraints obtained on NP in the $B_q$ system by including this additional information.

In the SM, $M^q_{12}$ is given by

$$M^q_{12}^{SM} = \frac{G_F^2 M^2_W}{12\pi^2} M_{B_q} \eta^B \bar{f}_{B_q} (V^*_{ts} V_{tb})^2 S_0(x_t),$$

where $G_F$ is Fermi’s constant, $M_W$ the mass of the $W$ boson, $\eta^B = 0.551$ a short-distance QCD correction (which is the same for the $B^0_q$ and $\bar{B}^0_q$ systems) $^7$, whereas the bag parameter $\bar{f}_{B_q}$ and the decay constant $f_{B_q}$ are non-perturbative quantities. $V_{ts}$ and $V_{tb}$ are elements of the Cabibbo–Kobayashi–Maskawa (CKM) matrix, and $S_0(x_t) \equiv \overline{m}_{t}/M_W$ = 2.32 $\pm$ 0.04 with $m_t(m_t) = (163.4 \pm 1.7)$ GeV, Ref. $^8$, describes the $t$-quark mass dependence of the box diagram with internal $t$-quark exchange; the contributions of internal $c$ and $\nu$ quarks are suppressed by $(m_{c,\nu}/M_W)^2$, by virtue of the GIM mechanism. Thanks to the suppression of light-quark loops, $M^q_{12}$ is dominated by short-distance processes and sensitive to NP.

In the SM, the mixing phase in the $B_s$ system is given by

$$\phi^S_{B_s} = -2\lambda^2 R_b \sin \gamma \approx -2^\circ,$$

where $\gamma$ is one of the angles of the unitarity triangle (UT), $\lambda$ is the Wolfenstein parameter and

$$R_b = \left(1 - \frac{\lambda^2}{2}\right) \frac{1}{\lambda} \frac{|V_{ub}|}{|V_{cb}|}.$$
Figure 1: Lines of constant $\rho_s$ (left) and constant $\phi_s^{\text{NP}}$ (right) in the $\sigma_s$-$\kappa_s$ plane. Blue line: $\rho_s \equiv 1$.

Up-to-date values of $\gamma$ from various sources can be found in Ref. 9, whereas $|V_{ub}|$ and $|V_{tb}|$ can be found in Refs. 10 and 11, respectively. The corresponding results from global fits can be found in Ref. 12.

In the presence of NP, the matrix element $M^q_{12}$ can be written, in a model-independent way, as

$$M^q_{12} = M^q_{12}^{\text{SM}} \left(1 + \kappa_q e^{i\phi_q}\right),$$

where the real parameter $\kappa_q \geq 0$ measures the “strength” of the NP contribution with respect to the SM, whereas $\sigma_q$ is a new CP-violating phase. Relating $\kappa_s$ and $\sigma_s$ to $\Delta M_s$, one has

$$\rho_s \equiv \frac{|\Delta M_s|}{\Delta M_s^{\text{SM}}} = \sqrt{1 + 2\kappa_s \cos \sigma_s + \kappa_s^2}. \quad (8)$$

The lines of $\rho_s = \text{const.}$ in the $\sigma_s$-$\kappa_s$ plane are shown in Fig. 1. The blue line $\rho_s = 1$ illustrates that even if the experimental value of $\Delta M_s$ coincides with the SM expectation, it is wrong to conclude that there is no NP in $B_s$ mixing - in fact, in this case the NP amplitude can be larger than the SM amplitude, i.e. $\kappa_s > 1$, if SM and NP contributions differ by a phase $\sigma_s$ between 120° and 240°.

In order to obtain $\rho_s$ from the experimental result (4), one has to determine $\Delta M_s^{\text{SM}}$. In addition to the input parameters listed after (5), one also needs the CKM matrix elements $|V_{ub}^* V_{tb}|$ and the hadronic matrix element $f^{B_s}_B$. The former is accurately known in terms of $|V_{cb}|$ and $\lambda$ and reads

$$|V_{ub}^* V_{tb}| = \left\{1 - \frac{1}{2} (1 - 2R_b \cos \gamma) \lambda^2 + O(\lambda^4)\right\} |V_{cb}^* V_{tb}| = (41.3 \pm 0.7) \times 10^{-3}. \quad (9)$$

The hadronic matrix element $f^{B_s}_B$ has been the subject of numerous lattice calculations, both quenched and unquenched, using various lattice actions and implementations of both heavy and light quarks. The current front runners are unquenched calculations with 2 and 3 dynamical quarks, respectively, and Wilson or staggered light quarks. Despite tremendous progress in recent years, the results still suffer from a variety of uncertainties which is important to keep in mind when interpreting and using lattice results. The most recent (unquenched) simulation by the JLQCD collaboration 13, with non-relativistic $b$ quarks and two flavours of dynamical light (Wilson) quarks, yields

$$f^{B_s}_B \left|_{\text{JLQCD}} = (0.245 \pm 0.021^{+0.003}_{-0.002}) \text{GeV}, \quad (10)$$

where the first error includes uncertainties from statistics and various systematics, whereas the second, asymmetric error comes from the chiral extrapolation from unphysically large light-quark masses to the $s$-quark mass.
More recently, (unquenched) simulations with three dynamical flavours have become possible using staggered quark actions. The HQQCD collaboration obtains\textsuperscript{14}

$$f_{B_s}B_{1/2}^{1/2}|_{HQQCD} = (0.281 \pm 0.021) \text{ GeV}.$$ \hspace{1cm} (11)

where all errors are added in quadrature.

Although we shall use both (10) and (11) in our analysis, we would like to stress that the errors are likely to be optimistic. There is the question of discretisation effects (LQCD uses data obtained at only one lattice spacing) and the renormalisation of matrix elements (for lattice actions without chiral symmetry, the axial vector current is not conserved and $f_{B_s}$ needs to be renormalised), which some argue should be done in a non-perturbative way\textsuperscript{15}. Simulations with staggered quarks also face potential problems with unitarity, locality and an odd number of flavours (see, for instance, Ref.\textsuperscript{15}). A confirmation of the HQQCD results by simulations using the (theoretically better understood) Wilson action with small quark masses will certainly be highly welcome.

With the above input parameters, one finds

$$\Delta M_s|_{LQCD} = (16.1 \pm 2.8) \text{ ps}^{-1}, \quad \Delta M_s|_{HQQCD} = (21.3 \pm 3.2) \text{ ps}^{-1},$$

$$\rho_s|_{LQCD} = 1.10 \pm 0.19, \quad \rho_s|_{HQQCD} = 0.83 \pm 0.13.$$ \hspace{1cm} (12)

The corresponding constraints in the $\sigma_s-\kappa_s$ plane are shown in Fig. 2.

In order to further constrain the NP parameter space, one needs to include information on the NP CP-violating phase $\phi_s^{NP}$. At the time Ref.\textsuperscript{2} was written, no such information was available. In the meantime, $\phi_s^{NP}$ has been constrained from measurements by the D0 collaboration, of the CP-asymmetry in flavour-specific (semileptonic) $B_s$ decays\textsuperscript{3} and the time-dependent angular analysis of untagged $B_s \to J/\psi \phi$ decays\textsuperscript{4}. These measurements can be translated, using supplementary information on the semileptonic asymmetry in $B_d$ decays, in the following results for $\Delta \Gamma_s$ and $\phi_s$:\textsuperscript{5,6}

$$\Delta \Gamma_s = \Gamma_L - \Gamma_H = (0.13 \pm 0.09) \text{ ps}^{-1}, \quad \phi_s = -0.70^{+0.47}_{-0.30}.$$ \hspace{1cm} (13)

These results actually are determined only up to a 4-fold ambiguity for $\phi_s$ and the sign of $\Delta \Gamma_s$: $\phi_s \to \pm \phi_s$ for $\Delta \Gamma_s > 0$ and $\phi_s \to \pm (\pi - \phi_s)$ for $\Delta \Gamma_s < 0$. As the SM prediction for $\phi_s$ is close to $0^\circ$, we can identify this result with $\phi_s^{NP}$. The combined constraints posed by $\Delta M_s$ and $\phi_s^{NP}$

\textsuperscript{5}Which also implies that we do not have to distinguish our definition of $\phi_s$ as $\text{arg}M_{12}$ from the definition used by the D0 collaboration, $\phi_s = \text{arg}(M_{12}/M_{12}^*)$.\textsuperscript{6}
on the new-physics parameters $\kappa_s$, $\sigma_s$ are shown as green areas in Fig. 2, including the 4-fold ambiguity. It is evident that at present the experimental error of $\phi_s^{NP}$ is too large to considerably reduce the area constrained by $\Delta M_s$ alone. The ambiguity can be reduced to a 2-fold one if some theory-input about the signs of $\cos\delta_{t2}$ is used, where $\delta_{t2}$ are the strong phases involved in the angular analysis of $B_s \rightarrow J/\psi\phi$, Ref.\textsuperscript{17}. At LHCb, it will be possible to study the time-dependence of flavour-tagged $B_s$ decays, which gives access to the mixing-induced asymmetry and allows one to reduce the number of discrete ambiguities without input from theory.

The above results can be compared with the following recent theory prediction for $\Delta \Gamma_s$, which is based on an improved operator product expansion of $\Gamma_{t2}$, the off-diagonal element of the $B_s$ decay matrix:\textsuperscript{18}

$$\Delta \Gamma_s^{\text{th}} = (0.096 \pm 0.039) \text{ ps}^{-1},$$

which agrees with the experimental result (13) within errors. Ref.\textsuperscript{18} also contains a detailed discussion of the theoretical predictions for flavour-specific CP asymmetries both in the $B_d$ and $B_s$ system and the constraints on NP in $B_s$ mixing extracted from all available experimental data.

Let us conclude with a few remarks concerning the prospects for the search for NP through $B^0_s\bar{B}^0_s$ mixing at the LHC. This task will be very challenging if essentially no CP-violating effects will be found in $B_s \rightarrow J/\psi\phi$ (and similar decays). On the other hand, even a small phase $\phi_s^{NP} \approx -10^\circ$ would lead to CP asymmetries at the $-20\%$ level, which could be unambiguously detected after a few years of data taking, and would not be affected by hadronic uncertainties. Ref.\textsuperscript{19} quotes a sensitivity to $\phi_s$ of $\sigma(\phi_s) = 1.2^\circ$ for an integrated luminosity of $2 \text{fb}^{-1}$ at LHCb and a sensitivity $\sigma(\Delta \Gamma_s/\Gamma_s) \sim 0.01$ for both LHCb and Atlas/CMS (at $30 \text{fb}^{-1}$). Conversely, the measurement of such an asymmetry would allow one to establish a lower bound on the strength of the NP contribution – even if hadronic uncertainties still preclude a direct extraction of this contribution from $\Delta M_s$ – and to dramatically reduce the allowed region in the NP parameter space. In fact, the situation may be even more promising, as specific scenarios of NP still allow large new phases in $B^0_s\bar{B}^0_s$ mixing, also after the measurement of $\Delta \Gamma_s$, see, for instance, Refs.\textsuperscript{20,21}.

In essence, the lesson to be learnt from this discussion is that NP may actually be hiding in $B^0_s\bar{B}^0_s$ mixing, but is still obscured by parameter uncertainties, some of which will be reduced by improved statistics at the LHC, whereas others require dedicated work of, in particular, lattice theorists. The smoking gun for the presence of NP in $B^0_s\bar{B}^0_s$ mixing will be the detection of a non-vanishing value of $\phi_s^{NP}$ through CP violation in $B_s \rightarrow J/\psi\phi$. This example is yet another demonstration that flavour physics is not an optional extra, but an indispensable ingredient in the pursuit of NP, also and in particular in the era of the LHC.

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    N. Magini (for Atlas and CMS), talk given at CKM06.
Tevatron results on b-hadron lifetimes and rare decays

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We present the most recent measurements on b-hadron lifetimes and rare decays performed by the CDF and DØ Collaborations at the Tevatron Collider, FERMILAB. We report the first Tevatron measurement based on 2 fb\(^{-1}\) of data collected during Run II.

Keywords: Tevatron, lifetimes, rare decays

1 Introduction

Since the start of RunII at the Tevatron \(pp\) Collider at the Fermi National Accelerator Laboratory (Fermilab), CDF and DØ experiments have produced a substantial amount of b–physics results. The Tevatron b–physics program is complementary to the program of the B factories, where a clear understanding of \(B^0\) and \(B^+\) mesons has been achieved. Currently, heavier b–hadron production is only possible at the Tevatron. Although the environment at \(pp\) is not as clean as in the \(e^+e^-\) case, it has the advantage of the high \(b\bar{b}\) production cross section and the high integrated luminosity.

In this document, we present the most recent results obtained by the Tevatron experiments: CDF and DØ. In Sect. 2, we discuss the new and most precise lifetime measurements for b–hadrons; \(B_s\), and \(\Lambda_b\); using data samples of integrated luminosity ranging from 0.4 to 1.3 fb\(^{-1}\). In Sect. 3, we show the most precise limits on \(B_{s,d} \rightarrow \mu^+\mu^-\) rare decays. We will also present the first ever measurement made at the Tevatron from an integrated luminosity dataset of 2 fb\(^{-1}\).
2 Lifetimes

Over the last years, the understanding of b-hadron lifetimes has improved considerably. Models, like the Heavy Quark Expansion (HQE)\(^1\), allow for a systematic expansion in orders of \(\alpha_s\) and \(1/\text{mass}^2\) of the total decay widths of heavy-quark hadrons. HQE gives very precise predictions on the ratios of the different b-hadron lifetimes. All these predictions have been in good agreement with the experimental data, except for the \(\Lambda_b\) baryon. Here the measured lifetime ratio \(\tau(\Lambda_b)/\tau(B_d)\) is significantly below the theoretical prediction. This has caused a review of all theoretical predictions. Recently improved lattice gauge theory computations for b-hadron lifetimes\(^2\) have decreased this theoretical expectation substantially to 0.88 ± 0.05, which now is in better agreement with the experimental data.

Both experiments at the Tevatron, CDF and DØ, have measured the lifetime of \(\Lambda_b\) using the exclusive decay channel \(\Lambda_b \rightarrow J/\psi(\rightarrow \mu^+\mu^-)\Lambda(\rightarrow p\pi)\). For purposes of calibration and normalization they used the \(B_d\) decay channel \(B_d \rightarrow J/\psi K_S(\rightarrow \pi^+\pi^-)\), which has a very similar final state topology to the \(\Lambda_b\) decay. Apart from that, the lifetime of the \(B_d\) is very well known from the \(B\) factories. When reconstructing these two final states the \(J/\psi\) decay vertex is combined with the reconstructed \(\Lambda\) (or \(K_S\)) track to form the \(\Lambda_b\) (\(B_d\)) vertex in the plane perpendicular to the beam direction; correcting for the boost gives the proper decay length (PDL) estimate. Then, a simultaneous unbinned maximum likelihood fit of the invariant mass and PDL distributions (taking into account the event-by-event PDL resolution) is performed to extract the lifetime.

![Figure 1: CDF. Likelihood fit projection in the proper decay length distribution for \(\Lambda_b \rightarrow J/\psi \Lambda^0\) decay channel in the signal region. Dashed line represents the \(\Lambda_b\) signal.](image)

CDF and DØ have measured the \(B_d\) lifetime compatible with the world average. While, for the \(\Lambda_b\), DØ has measured a lifetime ratio \(\tau(\Lambda_b)/\tau(B_d) = 0.811^{+0.036}_{-0.035}\) (stat.) ± 0.034 (syst.)\(^4\) and CDF \(\tau(\Lambda_b)/\tau(B_d) = 1.018 ± 0.062\) (stat.) ± 0.007 (syst.)\(^3\) they are different but still compatible with previous measurements and the theoretical predictions. Fig. 1 shows the likelihood fit projection in the proper decay length distribution for CDF \(\Lambda_b\) lifetime measurement, while the proper decay length distribution and the result of the fit for DØ measurement is shown in the left plot of figure 2. CDF result is higher than previous measurements, which are dominated by semileptonic decays.

DØ has recently reported a measurement of the \(\Lambda_b\) lifetime using the inclusive semileptonic decay channel \(\Lambda_b \rightarrow \Lambda^+ \mu^- \nu\). The measurement benefits from large statistics of semileptonic decay channels, but suffers from not being able to observe the actual \(\Lambda_b\) "peak". DØ reconstructs the \(\Lambda^+_c\) baryon using its decay channel \(\Lambda^+_c \rightarrow K_S p\). This particle is observed on top of a large
background, which is subtracted statistically, in bins of the visible proper decay length. All sources of remaining background to the \( \Lambda_b \) decay as well as corrections for the boost are properly taken into account by means of Monte Carlo simulations. The measured lifetime, \( \tau(\Lambda_b) = 1.28^{+0.12}_{-0.11} \) \((\text{stat.)} \pm 0.09(\text{syst.}) \) ps \(^4\), is in a good agreement with the measurement performed by DØ in the \( J/\psi \Lambda \) decay channel. The measured yields in the virtual proper decay length bins and the result of the lifetime fit are shown in the right plot of figure 2. As we can see, even when the apparent differences between theory and experiment have been resolved we still have an experimental issue to address to fully understand the \( \Lambda_b \) lifetime.

Tevatron experiments have also measured the lifetimes of the other \( b \)-hadrons. Using a large \( J/\psi \) sample, CDF has measured the lifetimes of \( B^+ \) and \( B_s \) mesons through theirs fully reconstructed decay channels. For \( B^+ \), by reconstructing the decay channel \( B^+ \rightarrow J/\psi K^+ \) they have measured \( \tau(B^+) = 1.630 \pm 0.016(\text{stat.}) \pm 0.011(\text{syst.}) \). For \( B_s \) they reconstructed \( B_s \rightarrow J/\psi \phi \), follow by \( \phi \rightarrow K^+K^- \), and have measured a lifetime of \( \tau(B_s) = 1.494 \pm 0.054(\text{stat.}) \pm 0.009(\text{syst.}) \) ps \(^3\). Both measurements are in good agreement with earlier measurements, as well as with HQE predictions. DØ has measured the \( B_s \) lifetime using the flavor specific decay channel \( B_s \rightarrow D^-_{s} \mu^+ \nu_X \), with the \( D^-_{s} \) reconstructed through \( D^-_{s} \rightarrow \phi \tau^- \), and \( \phi \rightarrow K^+K^- \). They find \( \tau(B^+_s) = 1.398 \pm 0.044(\text{stat.})^{+0.026}_{-0.028}(\text{syst.}) \) ps \(^3\). Fig. 3 shows the pseudo proper decay length distribution for \( D^-_{s} \mu^+ \) candidates with the result of the fit superimposed. This measurement is the current best world measurement, even more precise than the current world average. The measurement is also in fair agreement with previous measurements. It is important to notice that current lifetime measurements are reaching precision of the order 1-3% in the systematic uncertainty determination and with further data accumulation more incisive tests of HQE can be achieved.

3 Rare decays

Searches for the decays \( B_s \rightarrow \mu^+\mu^- \) and \( B_d \rightarrow \mu^+\mu^- \) can be a powerful tool to probe for physics beyond the Standard Model (SM). The purely leptonic decay \( B_{d,s} \rightarrow \mu^+\mu^- \) is a Flavor-Changing Neutral Current (FCNC) process. In the SM, this decay is forbidden at the tree level and even when higher order contributions are taken into account the predicted rate is very low. The SM branching ratio for this channel was first calculated \(^6\) and later refined to include QCD corrections \(^7\). The latest SM prediction \(^8\) is, \( Br(B_{s,d} \rightarrow \mu^+\mu^-) = (3.42\pm0.54) \times 10^{-9} \), where the error is dominated by non-perturbative hadronic uncertainties. The corresponding leptonic
branching fraction for the $B_d$ meson is suppressed by an additional factor of $|V_{td}/V_{ts}|^2$ leading to an expected SM branching ratio of $(1.00\pm0.14)\times10^{-10}$. There are various extensions to the SM that predict an enhancement of this branching ratio by 1 to 3 orders of magnitude.

CDF and DØ have performed very sophisticated likelihood ratio studies of $B_{s,d} \rightarrow \mu^+\mu^-$ starting with dimuon trigger data samples. Using kinematic and topological information based on Monte Carlo predictions they have been able to isolate a tiny set of candidate events. Then using the $B^+ \rightarrow J/\psi(\rightarrow \mu^+\mu^-)K^+$ decays as normalization (where the efficiency of the reconstructed two muons almost cancel in the ratio) they are able to extract the branching ratio limit for $B_{s,d} \rightarrow \mu^+\mu^-$. CDF searches using $0.78\ fb^{-1}$ of data have found 1 candidate event for $B_s$ and 2 candidates for $B_d$ yielding 90% C.L. limits on the branching ratio of $8.0 \times 10^{-8}$ and $2.3 \times 10^{-8}$, respectively. Fig. 4 shows the invariant mass distribution of surviving events versus the likelihood ratio distribution. Meanwhile DØ using $2\ fb^{-1}$ of accumulated data has found 3 candidate events for $B_s$ which translate to a limit on the branching ratio at the 90% C.L. of $7.5 \times 10^{-8}$. Fig. 5 shows the invariant mass distribution of surviving events for RunIIa and RunIIb (before and after the last summer shutdown). These measurements are the best existing experimental bounds and are just 20 times larger than the SM prediction for the case of $B_s$ and 300 times for the case of $B_d$. Improving analysis techniques and accessing data as they become available will allow us to reach limits close to the SM expectations.

4 Summary

We have presented the most recent new measurements of $b$–hadron lifetimes and rare decays from the Tevatron, in all cases an improve understanding of the the experimental uncertainty is evident. CDF has reached experimental uncertainties on the lifetime determination of an order of 1% in fully reconstructed decay channel. DØ has reached 2% level for the case of semileptonic decay channels. Small discrepancies between the DØ and CDF measurements are observed, further studies should shed light on them. The limits on $B_s \rightarrow \mu^+\mu^-$ are getting closer to the SM prediction, further data accumulation will soon rule out or confirm the known extensions of the SM.
Figure 4: CDF. Dimuon candidates invariant mass versus likelihood ratio distribution with final selection.

Figure 5: DØ. Dimuon invariant mass distribution of survival candidate events.

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Determination of $\phi_2$ ($\alpha$) and $\phi_3$ ($\gamma$)

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We review recent measurements constraining the Cabibbo-Kobayashi-Maskawa weak phases $\phi_2$ ($\alpha$) and $\phi_3$ ($\gamma$) at two energy-asymmetric $B$-factories, Belle/KEKB and BaBar/PEP-II. The measurements of decay processes $B \rightarrow \pi\pi$, $\rho\pi$, and $\rho\pi$ yield dominant constraint on $\phi_2$. The recent measurement of the process $B \rightarrow \alpha_3\pi$ is also sensitive to $\phi_2$. The phase $\phi_3$ is mainly measured using the decay process of $B^+ \rightarrow D^{*+}K^+$.

1 Introduction

$CP$ violation arises from irreducible phase in Cabibbo-Kobayashi-Maskawa (CKM) matrix$^{1,2}$ which is established as a part of the standard model (SM) by various experimental observations, in particular those from the asymmetric energy $B$-factories.$^{3,4}$ Now, we have entered a new era of the physics of $CP$ violation; we use it as a prove sensitive to the physics beyond the SM, or new physics. To examine whether all the observables are universally described by the CKM picture is a quest for the new physics, since we expect to see deviations from the CKM expectations if there contributes the new physics effect.

The test of the unitarity of the CKM matrix, $V_{ud}V_{ub}^*+V_{cd}V_{cb}^*+V_{td}V_{tb}^*=0$, is one of the most important new physics searches in this sense. The angles $\phi_2$ and $\phi_3$ denote the $CP$-violating weak phases $\arg[-V_{ub}V_{ub}^*/V_{ud}V_{ub}]$ and $\arg[-V_{ud}V_{ub}^*/V_{cd}V_{ub}]$, respectively, and their measurements are crucial for the test of the unitarity.

In this proceedings, we present the review of recent measurements of the angles $\phi_2$ and $\phi_3$ at Belle/KEKB and BaBar/PEP-II.

*Another naming convention, $\alpha(=\phi_2)$ and $\gamma(=\phi_3)$, are also used in the literature.
2 Determination of $\phi_2 (\alpha)$

The angle $\phi_2$ is the complex phase of $-V_{td}V_{tb}^*/V_{ud}V_{ub}^*$ and we can access it via the decay processes that involve those CKM factors. As shown in Fig. 1, the factor $V_{td}V_{tb}^*$ arises from $B^0$-$\bar{B}^0$ mixing and $V_{ud}V_{ub}^*$ appears in the decay processes that involve $b \to u$ transition, such as $B^0 \to \pi^+\pi^-$, $\rho^+\rho^-$, $\rho^+\pi^+$, and $a_1^+\pi^+$; we can constrain $\phi_2$ by the measurements of these processes.

In the decay chain $\Upsilon(4S) \to B^0\bar{B}^0 \to ff_{\text{tag}}$, where $f$ is the final states of interest (e.g., $\pi^+\pi^-$, $\rho^+\rho^-$, $\rho^+\pi^+$, and $a_1^+\pi^+$) and $f_{\text{tag}}$ is a final state that distinguishes $B^0$ and $\bar{B}^0$, the time-dependent differential decay rate is

$$\frac{df}{d\Delta t} \propto e^{-|\Delta t|/\tau_B} \left[ 1 + q_{\text{tag}} \left( A_f \cos(\Delta m_d \Delta t) + S_f \sin(\Delta m_d \Delta t) \right) \right]. \quad (1)$$

Here, $q_{\text{tag}}$ is the $b$-flavor charge [$q_{\text{tag}} = +1(-1)$ when $f_{\text{tag}}$ is a $B^0$ ($\bar{B}^0$) flavor eigenstate], $\Delta t$ is the decay time difference between the two $B$ mesons ($t_f - t_{f_{\text{tag}}}$), and $\tau_B$ and $\Delta m_d$ are the average lifetime and mass difference of the mass eigenstates of neutral $B$ mesons, respectively. The $CP$-violation parameters $A_f$ and $S_f$ are the observables of the measurement. In a decay process of $b \to u$ transition, $S_f$ is related to $\phi_2$ as

$$S_f = \sin 2\phi_2, \quad (2)$$

ignoring the penguin diagram contribution described in the followings.

The difficulty in the measurement of $\phi_2$ through the decays with $b \to u$ transition arises out of the possible contribution from the gluonic penguin diagram of $b \to d$ transition (Fig. 2). This contribution contaminates the measurement of $\phi_2$, since the decays via this diagram yields the CKM factor $V_{td}V_{tb}$, which is different from that of the $b \to u$ transition. With the contamination, Eq. (2) becomes

$$S_f = \sqrt{1 - A_f^2} \sin 2\phi_2^{\text{eff}}, \quad (3)$$

where $\phi_2^{\text{eff}}$ may differ from $\phi_2$. Large amount of effort (both experimental and theoretical) is devoted to constrain the difference between $\phi_2^{\text{eff}}$ and $\phi_2$, as well as the measurement of the $CP$-violation parameters $A_f$ and $S_f$.

The isospin analysis\(^8\) gives constraint on the difference $|\phi_2^{\text{eff}} - \phi_2|$ by incorporating the knowledge of processes related the modes. In the case of $B \to \pi\pi$, for example, it involves the decay modes $B^0 \to \pi^+\pi^-$ and $B^+ \to \pi^+\pi^0$ in addition to $B^0 \to \pi^+\pi^-$; one can constrain $|\phi_2^{\text{eff}} - \phi_2|$ by the measurements of branching fractions and $CP$ asymmetries of these modes. There are several methods other than the isospin analysis for the removal of the penguin diagram contamination. One is the time-dependent Dalitz plot analysis in $B^0 \to (\rho\pi)^0 \to \pi^+\pi^-\pi^0$ decay process,\(^9\) which is used along with an isospin relation.\(^10,11\) Another is the use of theoretical assumptions,\(^12,13,14,15\) such as the flavor $SU(3)$ symmetry; it enables us to relate the $CP$-violation measurement of $B^0 \to a_1^+\pi^+$ to $\phi_2$, as well as the decay modes of $B^0 \to \pi^+\pi^-$, $\rho^+\rho^-$, and $\rho^+\pi^+$.

![Feynman diagrams](image-url)
2.1 $B \to \pi \pi$

Belle and BaBar measure the $CP$ violation in the decay process $B^0 \to \pi^+\pi^-$ with data samples containing $535 \times 10^6$ and $383 \times 10^6$ $B\bar{B}$ pairs, respectively.\textsuperscript{16,17} Table 1 lists the results. Belle reports the first observation of the direct $CP$ violation ($A_{\pi^+\pi^-} \neq 0$) with a significance of 5.5 standard deviations ($\sigma$). BaBar observes $CP$ violation ($C_{\pi^+\pi^-} \neq 0$ and $S_{\pi^+\pi^-} \neq 0$) with a significance of 5.4 $\sigma$, though it does not see significant direct $CP$ violation. The discrepancy between the results from Belle and BaBar corresponds to the significance of 2.1 $\sigma$. From the above result and world averaged values for $B(B^0 \to \pi^+\pi^-), B(B^0 \to \pi^0\pi^0), B(B^+ \to \pi^+\pi^0)$, and $A(B^0 \to \pi^0\pi^0)$,\textsuperscript{18} Belle constrain $\phi_2$ and excludes the region $11^\circ < \phi_2 < 79^\circ$ at the 95% confidence level (C.L.). Further discussion on the $\phi_2$ constraint using the above mentioned results can be found elsewhere.\textsuperscript{19,20}

Table 1: The results of the latest measurements of $CP$ violation in the decay $B^0 \to \pi^+\pi^-$ at Belle and BaBar. The first and second errors are statistical and systematic, respectively. Note that they use different conventions for the coefficient of $\cos(\Delta m_t\Delta t)$, $A_{\pi^+\pi^-}$ and $C_{\pi^+\pi^-}$, which are equivalent except for the opposite sign.

<table>
<thead>
<tr>
<th></th>
<th>$A_{\pi^+\pi^-}$</th>
<th>$S_{\pi^+\pi^-}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Belle</td>
<td>+0.55 ± 0.08 ± 0.05</td>
<td>-0.61 ± 0.10 ± 0.04</td>
</tr>
<tr>
<td>BaBar</td>
<td>+0.21 ± 0.09 ± 0.02</td>
<td>-0.60 ± 0.11 ± 0.03</td>
</tr>
</tbody>
</table>

2.2 $B \to \rho \rho$

There are two essential differences between $B \to \rho \rho$ and $B \to \pi \pi$: 1) $\rho$ is a vector meson, while $\pi$ is a pseudo-scalar meson, and 2) the contribution from the $b \to d$ penguin diagram is small in $B \to \rho \rho$ than in $B \to \pi \pi$. In principle, the first difference is a disadvantage of $B \to \rho \rho$, since it leads to the fact that the final state is a mixture of three polarization states with different $CP$ eigenvalues. In practice, however, it is not because one of the polarization state, the longitudinal polarization, turned out to dominate the final state in all of the processes of $B^0 \to \rho^+\rho^-, \rho^0\rho^0$, and $B^+ \to \rho^+\rho^0$.\textsuperscript{21,22,23,34,35}; this enables us to relate the measurements to $\phi_2$ without the effort to disentangle a single polarization state. The second difference, the smallness of the penguin diagram contribution, is deduced by the smallness of the branching fraction of $B^0 \to \rho^0\rho^0$,\textsuperscript{23} compared to those of $B^0 \to \rho^+\rho^- + B^+ \to \rho^+\rho^0$. The smaller penguin leads to the smaller $|\phi_2^{\text{eff}}| < |\phi_2|$ and thus the better constraint on $\phi_2$.

Table 2 lists the latest measurements of the $CP$ violation in the process $B^0 \to \rho^+\rho^-$ at Belle and BaBar using data samples containing $535 \times 10^6$ and $347 \times 10^6$ $B\bar{B}$ pairs, respectively. The results are consistent with $S_{\rho^+\rho^-} = 0$, which correspond to $\phi_2 \sim 90^\circ$ at the limit of no penguin contribution. A drastic change in the $\phi_2$ constraint came from the new measurements of the branching fractions of $B^0 \to \rho^0\rho^0$ and $B^+ \to \rho^+\rho^0$.\textsuperscript{26,23} (Table 3) Before the new measurements,
the isospin triangles were in squashed shapes due to the small branching fractions of $B^0 \rightarrow \rho^+ \rho^-$ and $\rho^0 \rho^0$ compared to $B^+ \rightarrow \rho^+ \rho^0$, leading to the good constraint on $|\phi_2 - \phi_2^{\text{eff}}|$; with the new measurements, the triangles are now normally shaped and the constraint gets modest. For further improvement of the $\phi_2$ constraint, we have to measure the $CP$ asymmetry of $B^0 \rightarrow \rho^0 \rho^0$ decay process or to assume some theoretical input such as flavor $SU(3)$.

From the above mentioned measurements, Belle obtains the constraint of $54^\circ < \phi_2 < 113^\circ$ at the 90% C.L., while BaBar obtains $\phi_2^{\text{eff}} = (95.5^{+6.9}_{-6.2})^\circ$ and $|\phi_2 - \phi_2^{\text{eff}}| < 18^\circ$ at 68% C.L. Further discussion can be found elsewhere.\textsuperscript{19,20}

Table 2: The latest measurements of the $CP$ violation, polarization, and branching fraction of the decay $B^0 \rightarrow \rho^+ \rho^-$ at Belle and BaBar. The first and second errors are statistical and systematic, respectively. Note that the branching fraction and polarization measurements at Belle uses a different data sample from that used for the $CP$ violation measurement.

<table>
<thead>
<tr>
<th></th>
<th>$A_{\rho^+ \rho^-}$ ($= -C_{\rho^+ \rho^-}$)</th>
<th>$S_{\rho^+ \rho^-}$</th>
<th>$B \times 10^{-6}$</th>
<th>$f_L$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Belle</td>
<td>$+0.16 \pm 0.21 \pm 0.07$</td>
<td>$+0.19 \pm 0.30 \pm 0.07$</td>
<td>$22.8 \pm 3.8^{+12.0}_{-6.6}$</td>
<td>$0.941^{+0.044}_{-0.040} \pm 0.030$</td>
</tr>
<tr>
<td>BaBar</td>
<td>$+0.07 \pm 0.15 \pm 0.06$</td>
<td>$-0.19 \pm 0.21^{+0.05}_{-0.07}$</td>
<td>$23.5 \pm 2.2 \pm 4.1$</td>
<td>$0.977 \pm 0.024^{+0.015}_{-0.013}$</td>
</tr>
</tbody>
</table>

Table 3: The latest measurements of the branching fractions and longitudinal fractions of the decays $B^0 \rightarrow \rho^0 \rho^0$ and $B^+ \rightarrow \rho^+ \rho^0$. The first and second errors are statistical and systematic, respectively.

<table>
<thead>
<tr>
<th></th>
<th>$B \times 10^{-6}$</th>
<th>$f_L$</th>
<th>$B \times 10^{-6}$</th>
<th>$f_L$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Belle</td>
<td>$31.7 \pm 7.1^{+3.5}_{-6.7}$</td>
<td>$0.948 \pm 0.106 \pm 0.021$</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>BaBar</td>
<td>$16.8 \pm 2.2 \pm 2.3$</td>
<td>$0.905 \pm 0.042^{+0.023}_{-0.027}$</td>
<td>$1.07 \pm 0.33 \pm 0.19$</td>
<td>$0.87 \pm 0.13 \pm 0.04$</td>
</tr>
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</table>

2.3 $B \rightarrow \rho \pi$

The time-dependent Dalitz plot analysis of the decay $B^0 \rightarrow \rho \pi \rightarrow \pi^+ \pi^- \pi^0$ offers a unique way to determine the angle $\phi_2$; in contrast to the other analyses, it includes measurement of $CP$-violating asymmetries in mixed final states. In the limit of high statistics, this enables us to obtain $\phi_2$ without discrete ambiguities.

Belle and BaBar have recently released the complete time-dependent Dalitz plot analysis of the $B^0 \rightarrow \pi^+ \pi^- \pi^0$ decay process for the first time using data samples corresponding to $449 \times 10^6$ and $375 \times 10^6 BB$ pairs, respectively.\textsuperscript{27,28} From the result, BaBar obtains the constraint $\phi_2 = (87^{+145}_{-133})^\circ$ at 68% C.L. Belle performs a combined analysis of the time-dependent Dalitz plot and isospin (pentagon) analysis and obtains $68^\circ < \phi_2 < 95^\circ$ as a 68% confidence interval consistent with the SM, though a large SM-disfavored region also remains. BaBar have also released the new measurement of the branching fraction $(10.2 \pm 1.4(\text{stat}) \pm 0.9(\text{syst}) \times 10^{-6})$ and $CP$ asymmetry $(A = -0.01 \pm 0.13(\text{stat}) \pm 0.02(\text{syst}))$ in the decay $B^+ \rightarrow \rho^+ \pi^0$,\textsuperscript{29} which can be used to improve the $\phi_2$ constraint from the pentagon analysis.

2.4 $B \rightarrow a_1 \pi$

The process $B^0 \rightarrow a_1^+ \pi^+$ is also described by the diagrams of Figs. 1 and 2 and one can obtain information on $\phi_2$ through the time-dependent $CP$ violation measurement of this process, though the final states $a_1^+ \pi^+$ are not $CP$ eigenstates. BaBar reports the measurement using a data sample containing $384 \times 10^6 BB$ pairs, obtaining $\phi_2^{\text{eff}} = 78.6^\circ \pm 7.3^\circ$, where the error includes both statistical and systematic errors.\textsuperscript{30} Since the isospin analysis is difficult to perform with this
process, one basically have to rely on theoretical assumption such as flavor $SU(3)$\textsuperscript{14} to extract $\phi_2$ from the experimentally measured $\phi_2^{\text{eff}}$.

3 Determination of $\phi_3$ ($\gamma$)

The angle $\phi_3$ is the complex phase of $-V_{ud}V_{ub}^*/V_{cd}V_{cb}^*$. As shown in Fig. 3, the factors $V_{ub}^*$ and $V_{us}(\sim V_{cd})$ appears in the process of $B^- \rightarrow D^{(*)0}K^-$ decay, while the process $B^- \rightarrow \overline{D}^{(*)0}K^-$ involves the factors $V_{ub}^*$ and $V_{cs}(\sim V_{ud})$. Thus, one can access $\phi_3$ via a decay process of $B^- \rightarrow fK^-$ where $f$ is a final state to which both $D^{(*)0}$ and $\overline{D}^{(*)0}$ can decay.

The decay amplitudes of the processes $B^+ \rightarrow DK^+(D \rightarrow \overline{f})$ and $B^- \rightarrow DK^-(D \rightarrow f)$ are written as

$$A(B^+ \rightarrow DK^+) \propto A_D + r_B e^{i\phi_3 + i\delta} A_{\overline{D}}$$

$$A(B^- \rightarrow DK^-) \propto A_D + r_B e^{-i\phi_3 + i\delta} A_{\overline{D}}$$

(4)

where $A_D$ and $A_{\overline{D}}$ are the decay amplitudes of $D^{(*)0} \rightarrow f$ and $\overline{D}^{(*)0} \rightarrow f$, respectively; $r_B = |A(B^- \rightarrow \overline{D}^{(*)0}K^-)|/|A(B^- \rightarrow D^{(*)0}K^-)|$; and $\delta$ is the strong phase difference between $B^- \rightarrow \overline{D}^{(*)0}K^-$ and $B^- \rightarrow D^{(*)0}K^-$. As can be seen in the equations, the dependence on $\phi_3$ can manifest itself as the interference between the first and second terms in the right hand side; the comparison of the interference effects in $B^+$ and $B^-$ decays is sensitive to $\phi_3$. The difficulty in using this decay process comes from the smallness of $r_B(\sim 0.1)$, which is mainly due to color suppression. The methods of the analysis can be categorized into the following three by the types of the final state $f$: Dalitz plot,\textsuperscript{31,32} GLW (Gronau-London-Wyler),\textsuperscript{33,34} and ADS (Atwood-Dunietz-Soni).\textsuperscript{35,36}

Another possibility to access $\phi_3$ is the time dependent $CP$ violation measurement of $B^0 \rightarrow D^{(*)\pm}\pi^\mp$ decay process, which is sensitive to $\sin(2\phi_1 + \phi_3)$.\textsuperscript{37,38} Since $\phi_1$ is very well measured compared to $\phi_3$, it is virtually a measurement of $\phi_3$.

![Feynman diagrams](image)

Figure 3: Feynman diagrams corresponding to the $B^- \rightarrow D^{(*)0}K^-$ decay (left) and $B^- \rightarrow \overline{D}^{(*)0}K^-$ (right).

3.1 Dalitz Plot Analysis

In the Dalitz plot analysis, one chooses the three body final state of $K_S\pi^+\pi^-$ as the common final state of $f$ and $\overline{f}$. In this case, the $D$ decay amplitudes of $A_D$ and $A_{\overline{D}}$ have Dalitz plot dependence as

$$A_D = f(m_{-}^2, m_{+}^2) \quad \text{and} \quad A_{\overline{D}} = f(m_{+}^2, m_{-}^2),$$

(5)

where $m_{-}$ and $m_{+}$ are the invariant mass of $K_S\pi^+$ and $K_S\pi^-$, respectively, and $f(m_{+}^2, m_{-}^2)$ is the complex Dalitz plot amplitude of $\overline{D}^0 \rightarrow K_S\pi^+\pi^-$ decay, which is well calibrated using large data samples of $e^+e^- \rightarrow \phi\phi$ events.

\textsuperscript{b}Here, no $CP$ violation in $D$ decays is assumed, i.e., we assume

$$A(D^{(*)0} \rightarrow f) = A(D^{(*)0} \rightarrow \overline{f}) \quad \text{and} \quad A(\overline{D}^{(*)0} \rightarrow f) = A(D^{(*)0} \rightarrow \overline{f}).$$
Both Belle and BaBar obtain the constraint on $\phi_3$ using this method.\textsuperscript{39,40} Belle analyzes the
decays of $B^+ \rightarrow DK^+$, $D^* K^+$, and $DK^{**}$ using a data sample containing $388 \times 10^6 \bar{B}B$ pairs
and obtains $\phi_3 = (53 \pm 15)_{\text{stat}} \pm 3_{\text{syst}} \pm 9_{\text{model}} \%$, while BaBar obtains $\phi_3 = (92 \pm 41)_{\text{stat}} \pm 11_{\text{syst}} \pm 12_{\text{model}} \%$ based on an analysis of the decays of $B^+ \rightarrow DK^+$ and $D^* K^+$ using a
data sample containing $347 \times 10^6 \bar{B}B$ pairs. Note that the significant difference of the statistical
errors between Belle and BaBar is due to their central values of $r_B$; BaBar happens to obtain
smaller $r_B$ than Belle, leading to the less sensitivity to $\phi_3$. BaBar has recently shown another
possibility of using the decay of $D \rightarrow \pi^+ \pi^- \pi^0$ for the Dalitz plot analysis.\textsuperscript{41}

3.2 GLW

In GLW, $CP$ eigenstates are chosen as the common final state $f$. With $CP$-even ($K^+K^-$,
$\pi^+\pi^-$) and $CP$-odd ($K_S\pi^0$, $K_S\omega$, $K_S\phi$, etc.) eigenstates denoted by $D_+$ and $D_-$, the direct
$CP$-violating asymmetry $A_\pm$ is calculated from Eq. (4) as

$$A_\pm = \frac{B(B^+ \rightarrow D_±K^-) - B(B^- \rightarrow D_±K^+)}{B(B^- \rightarrow D_±K^-) + B(B^+ \rightarrow D_±K^+)} = \frac{r_B \sin \delta \sin \phi_3}{1 + r_B^2 \pm 2r_B \sin \delta \sin \phi_3}.$$  \hspace{1cm} (6)

These are sensitive to $\phi_3$ combined with the additional information from double ratio of the
branching fractions

$$R_\pm = \frac{B(B \rightarrow D_±K)/B(B \rightarrow D_±\pi)}{B(B \rightarrow D^0 K)/B(B \rightarrow D^0 \pi)} = 1 + r_B^2 \pm 2r_B \sin \delta \sin \phi_3.$$  \hspace{1cm} (7)

Note that one has to observe significant direct $CP$ violation to constrain $\phi_3$ using this method
alone.

Belle and BaBar report the measurements with this method using data samples containing
$275 \times 10^6$ and $232 \times 10^6 \bar{B}B$ pairs, respectively.\textsuperscript{42,43,44} Though neither of the experiments observe
significant direct $CP$ violation, the information is still useful in the global fit to constrain $\phi_3$.

3.3 ADS

The ADS method uses the Cabibbo suppressed mode, such as $D^0 \rightarrow K^+\pi^-$, for the final state $f$. By this choice, the two decay processes of $B^- \rightarrow D^0 K^- (D^0 \rightarrow K^+\pi^-)$ and $B^- \rightarrow \bar{T}^0 K^- (\bar{T}^0 \rightarrow
K^+\pi^-)$ have comparably small branching fractions; the former is Cabibbo suppressed and the
latter is color suppressed. The ratio of the branching fractions between the favored and suppressed
decay processes is sensitive to $\phi_3$ as

$$R_{DK} = \frac{B^+ \rightarrow D(K^+\pi^-)K^-}{B^- \rightarrow D(K^-\pi^+)K^-} = r_B^2 + r_D^2 + 2r_B r_D \cos \phi_3 \cos \delta,$$  \hspace{1cm} (8)

where $r_D$ is the Cabibbo suppression factor defined as $r_D = |A(D^0 \rightarrow K^+ \pi^-)|/|A(\bar{T}^0 \rightarrow
K^+\pi^-)|$. The advantage of this method is that the interference effect is significant with the two
interfering decay processes having the similar branching fractions. The disadvantage is, on the
other hand, the smallness of the branching fraction of the suppressed decay process, leading to
large statistical uncertainty.

Belle and BaBar perform the analysis using data samples containing $275 \times 10^6$ and $232 \times
10^6 \bar{B}B$ pairs, respectively; neither of them see significant signal for the suppressed decay
processes.\textsuperscript{45,46} They obtain the constraint on $r_B$ based on the upper limit on the branching
fractions of the suppressed decays: at 90% C.L., Belle constrain $r_B < 0.18$ for the decay process
of $B \rightarrow DK$, while BaBar constrain $r_B < 0.23$ ($r_B < 0.16$) for the $B \rightarrow DK$ ($B \rightarrow D^*K$) decay
process. This knowledge can be used in the global fit to constrain $\phi_3$. 
3.4 Time-Dependent Analysis of $B^0 \to D^{(*)\pm}\pi^\mp$ for $\sin(2\phi_1 + \phi_3)$

The time-dependent decay rates of the $B^0 \to D^{(*)\pm}\pi^\mp$ decay processes are given by\textsuperscript{47}

$$P(B^0 \to D^{(*)\pm}\pi^\mp) \propto e^{-|\Delta t|/\tau_{D^*}} \left[ 1 \mp C \cos(\Delta m_d \Delta t) - S^\pm \sin(\Delta m_d \Delta t) \right], \quad (9)$$

$$P(\bar{B}^0 \to D^{(*)\pm}\pi^\mp) \propto e^{-|\Delta t|/\tau_{D^*}} \left[ 1 \pm C \cos(\Delta m_d \Delta t) + S^\pm \sin(\Delta m_d \Delta t) \right], \quad (10)$$

where $C$ and $S^\pm$ are the observables in the measurement. They can be related to $\sin(2\phi_1 + \phi_3)$ as

$$S^\pm = \frac{2(-1)^L R \sin(2\phi_1 + \phi_3 \pm \delta)}{1 + R^2}, \quad (11)$$

where $L$ is the angular momentum of the final state.

Belle and BaBar perform the time-dependent analysis using data samples containing $386 \times 10^6$ and $232 \times 10^6 B\bar{B}$ pairs, respectively.\textsuperscript{68,69,50} Belle obtains $|\sin(2\phi_1 + \phi_3)| > 0.44(0.13)$ and $|\sin(2\phi_1 + \phi_3)| > 0.52(0.07)$ at 68% (90%) C.L. using $B \to D^*\pi$ and $B \to D\pi$, respectively. BaBar constrains $\sin(2\phi_1 + \phi_3)$ combining $B \to D^{(*)}\pi$ and $B \to D\rho$ as $|\sin(2\phi_1 + \phi_3)| > 0.64(0.40)$ at 68% (90%) C.L.

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The Measurement of $\beta$ by Belle and BABAR

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We describe recent results on the measurement of the CP violation parameter $\beta$ from the Belle and BABAR experiments. Having established CP violation in $B$ decays, a number of distinct approaches have been utilized to search for deviations from the Standard Model expectation of CP violation in decays sensitive to $\sin 2\beta$ either at tree level or via penguin diagrams. New results from studies of $B \to (c\bar{c})K$, $B \to D^{(*)}\pi\pi$, $B \to D^{(*)}\pi\pi$ and $B \to \eta K$ decays include the first significant deviation from $\sin 2\beta = 0$ observed in charmless $B$ decays.

The presence of CP violation in $B$-meson decays is well established. Initial measurements of $B \to J/\psi K_S^0$ and related $B \to (c\bar{c})K$ decays by the Belle and BABAR experiments demonstrated that the Cabibbo Kobayashi Maskawa matrix within the Standard Model of electroweak physics provides the dominant mechanism for this CP violation.

In $B \to (c\bar{c})K$ decays, CP violation in the Standard Model arises from the interference of decay and mixing amplitudes. The measurement of this effect requires the determination of both the $B$ decay time difference, $\Delta t$, and flavor tag ($B^0$ or $\bar{B}^0$) of the decaying $B$ meson. While $B$ mesons are just below threshold at the $T(4S)$, the KEK and PEP-II accelerators utilize a boosted center of mass, which enables the measurement of $\Delta t$ with good resolution.

Flavor tagging relies upon distinctive features of $B^0$ decays, in particular the correlation of high momentum leptons as well as kaon charge with the flavor of the decaying $B$ meson. Excellent particle identification capabilities allow for efficient detection of electrons, muons, and kaons. Sophisticated algorithms developed by each experiment allow for an effective tagging efficiency of approximately 30%. Both vertex resolution and tagging performance are measured directly on the data.

Given sufficient $\Delta t$ resolution and flavor tagging capability, CP violation parameters are determined through a time-dependent analysis. For the decays studied here, the decay rate $f_+(f_-)$ when the tagging meson is a $B^0$ ($\bar{B}^0$) is given in terms of a complex parameter $\lambda$ by

$$f_+ (\Delta t) = \frac{e^{-|\Delta t|/\tau_{B^0}}}{4\Gamma_{B^0}} \left[ 1 \pm \frac{2Im \lambda}{1 + |\lambda|^2} \sin (\Delta m_d \Delta t) \pm \frac{1 - |\lambda|^2}{1 - |\lambda|^2} \cos (\Delta m_d \Delta t) \right],$$

(1)

where $\Delta t = t_{\text{rec}} - t_{\text{tag}}$ is the difference between the proper decay times of the reconstructed $B$ meson $B_{\text{rec}}$ and the tagging $B$ meson $B_{\text{tag}}$, $\tau_{B^0}$ is the $B^0$ lifetime, and $\Delta m_d$ is the $B^0$-$\bar{B}^0$ oscillation frequency. The CP asymmetry is then written

$$A_{CP}(\Delta t) \equiv \frac{f_+ - f_-}{f_+ + f_-} = S \sin (\Delta m_d \Delta t) - C \cos (\Delta m_d \Delta t),$$

(2)
where we have used the approximation that $\Delta t = 0$. In the modes described here, $S = -\eta_f \sin 2\beta$ and $C = 0$, where $\eta_f$ is the CP eigenvalue of the final state, in the Standard Model. We will quote results in terms of $\sin 2\beta_{eff} \equiv -\eta_f S$.

The large data samples collected by both the $B$-Factory experiments have subsequently allowed for the precise measurement of CP violation in $B \to (c\bar{c})K$ decays, as well as signs of physics beyond the Standard Model. Decay modes that are sensitive to the angle $\beta \equiv \arg [-V_{cb}V_{cb}^* / V_{td}V_{td}^* ]$ include $b \to c\bar{c}s$ decays, $b \to c\bar{c}d$ decays, and penguin-diagram dominated decays ($b \to d\bar{s}s$ and $b \to s\bar{s}s$).

In this contribution, we summarize a number of new results presented at this conference.

1 The Measurement of $\sin 2\beta$ with $B \to (c\bar{c})K$ decays

The most precise measurement of $\sin 2\beta$ is made using $B \to (c\bar{c})K$ decays, including the channels $J/\psi K_S^0$, $J/\psi K_L^0$, and for the BABAR analysis $\psi(2S)K_S^0$, $\chi_c K_S^0$, and $\eta_K K_S^0$. Most recent results from Belle and BABAR for these channels utilize approximately 535 million and 384 million $B\bar{B}$ events, respectively. The event samples from the Belle analysis are illustrated in Fig. 1. A total of nearly 25000 signal events are reconstructed in the two analysis. Modes with a $K_L^0$ are identified with a very high signal to noise. However, the $J/\psi K_L^0$ mode is more challenging: the $K_L^0$ is identified only via hadronic interactions in the calorimeter or detectors in the magnet flux return. Nevertheless, the $J/\psi K_L^0$ channel contributes substantially to the final CP result.

Figure 2 shows the $\Delta t$ distribution and raw asymmetry from the BABAR analysis. The characteristic $\sin(\Delta m_d \Delta t)$ distribution is apparent, as is the opposite CP asymmetry between the $\eta_f = -1$ and $\eta_f = +1$ channels. The results for $S$ are $^{2,3,4}$:

$$\sin 2\beta_{eff} = 0.714 \pm 0.032 \pm 0.018 \quad (\text{BABAR}),$$
$$\sin 2\beta_{eff} = 0.642 \pm 0.031 \pm 0.017 \quad (\text{Belle}),$$
$$\sin 2\beta_{eff} = 0.678 \pm 0.026 \quad (\text{Average}).$$

These measurements remain limited by their statistical precision, and the interpretation of these measurements in terms of $\sin 2\beta$ also does not introduce significant uncertainty $^5$. Also shown in Fig. 2 is a mode by mode breakdown of CP results presented for the first time by the BABAR experiment, including systematic uncertainties. No significant deviation from the average $\sin 2\beta_{eff}$ is observed. Additionally, both experiments fit for direct CP violation. No evidence is found, and the average result is $C = 0.002 \pm 0.021$. 

Figure 1: Event samples from the Belle analysis of $B \to J/\psi K_L^0$ (left) and $B \to J/\psi K_S^0$ (right). The points are the data, the sum of histograms shows the fit result, and the shaded histograms indicate the measured background level.
2 \ CP violation in $B \to D^+D^-$ and $B \to D^{*+}D^-$ decays

$CP$ violation in $B \to D^{(*)}D^{(*)}$ decays is related to $\sin2\beta$ if contributions from penguin diagrams are neglected. In the Standard Model, these contributions are expected to be only a few percent, so a significant deviation of $S$ ($C$) from $\sin2\beta$ (0) would be a strong indication of new physics contributions via penguin diagrams. Both Belle and BABAR have measured $B \to D^{*+}D^{*-}$ to have $CP$ violation in agreement with expectations.

Here, new results are summarized from both experiments on $B \to D^+D^-$, which has a lower branching fraction than $B \to D^{*+}D^{*-}$. The flavor-tagged $\Delta t$ distributions for each experiment are shown in Fig. 3. As is apparent, the experiments determine rather different results for $CP$ parameters based on their data samples.

The Belle experiment observes $CP$ violation at the 4.1σ level and determines

$$\begin{align*}
\sin2\beta_{eff} &= 1.13 \pm 0.37 \pm 0.09, \\
C &= -0.91 \pm 0.23 \pm 0.06,
\end{align*}$$

(4)

where the direct $CP$ coefficient is significantly different from expectations. This is evident from the asymmetry shown in number of events the $B^0$ and $\bar{B}^0$ $\Delta t$ distributions. The BABAR experiment instead determines

$$\begin{align*}
\sin2\beta_{eff} &= 0.54 \pm 0.34 \pm 0.06, \\
C &= 0.11 \pm 0.22 \pm 0.07,
\end{align*}$$

(5)

where $C$ is consistent with the Standard Model expectation (0). The Belle and BABAR result in approximately 3σ apart, thus further analysis is required to resolve the issue of direct $CP$-violation effects in $B \to D^+D^-$. In addition, BABAR has recently updated results on $B \to D^{*\pm}D^\mp$, which is not a $CP$ eigenstate. Here, The $D^{*+}D^-$ and $D^*-D^+$ decay modes are analyzed separately. In the absence of
Figure 3: Distributions of $B^0$ tagged (left) and $\bar{B}^0$ tagged $B \rightarrow D^+ D^-$ events from the Belle (top) and BABAR (bottom) analyses. The points show the data and the solid lines show the fit result.
Figure 4: Signal distributions for $B \to D^{*+}D^-$ (left) and $B \to D^{*-}D^+$ from BABAR. The points are the data and the curves show the total fit result as well as the background contribution.

$CP$ violation, $S_{D^{*+}D^-} = -S_{D^{*-}D^+}$ and $C_{D^{*+}D^-} = -C_{D^{*-}D^+}$. Figure 4 shows the signal samples for these modes, where BABAR determines

$$S_{D^{*+}D^-} = -0.79 \pm 0.21 \pm 0.06,$$
$$C_{D^{*+}D^-} = 0.18 \pm 0.15 \pm 0.06,$$
$$S_{D^{*-}D^+} = -0.44 \pm 0.22 \pm 0.04,$$
$$C_{D^{*-}D^+} = 0.23 \pm 0.15 \pm 0.04.$$  

(6)

Here, $S_{D^{*+}D^-} + S_{D^{*-}D^+}$ is significantly different from 0, a first indication of $CP$ violation in these decay modes.

3 Measurement of $CP$ violation in $B^0 \to D^0 \pi^0$ and related channels

$CP$ violation has been previously measured in $B^0 \to D^0 \pi^0$ and related color-suppressed channels utilizing the $D^0 \to K_0^0 \pi^+\pi^-$ final state and a Dalitz plot analysis. Here, new results are presented using other $CP$ eigenstate decays of the $D^0$ including $D^0 \to \pi^+\pi^-$, $D^0 \to K^+K^-$, and $D^0 \to K_0^0\omega$. Loop diagrams do not contribute to these decays, however $R$-parity-violating supersymmetric processes\(^9\) could enter at tree level in these decays. With a small interpretation uncertainty in the Standard Model, a significant deviation of $\sin 2\beta_{\text{eff}}$ from $\sin 2\beta$ in these decays would be a sign of new physics.

A total of 340 signal events are observed and the time-dependent analysis finds

$$\sin 2\beta_{\text{eff}} = 0.56 \pm 0.23 \pm 0.05$$
$$C = -0.23 \pm 0.16 \pm 0.04$$  

(7)

These results\(^10\) are in good agreement with Standard Model expectations. Figure 5 shows the $\Delta t$ distributions and raw asymmetries for the reconstructed $CP$ odd and even events.

4 Charmless $B$ decay measurements of $\sin 2\beta$

Unlike the $b \to c\bar{c}s$ amplitude dominated decays described above, $CP$ violation in decays to charmless final states proceed via a single loop (penguin) amplitude. For modes such as $B \to \eta'K$ and $B \to \phi K$, this amplitude has the same relative phase $\beta$ with respect to $B^0\bar{B}^0$ mixing and are thus sensitive to the same $CP$-violation parameter $\sin 2\beta$ in the Standard Model.
In these decays additional amplitudes, such as new physics processes contributing in the loop diagram, may give a sizable contribution to the observed CP-violation effect. A measurement of $\sin 2\beta$ in charmless decays that differs significantly from that in $b \to c\bar{c}s$ decays would be an indication of beyond the Standard Model physics. Theoretical estimates\,\textsuperscript{12} indicate that higher order Standard Model contributions are generally small and tend to increase the value of $\sin 2\beta_{\text{eff}}$ in these modes.

Analysis results from Belle and \textit{BABAR} indicate a possible deviation between $\sin 2\beta$ measured in $B \to (c\bar{c})K$ decays as described above and $b \to s\bar{s}s$ decays. Here we describe the first measurements\,\textsuperscript{4,11} of $\sin 2\beta_{\text{eff}}$ in charmless $B$ decays that conclusively differ from $\sin 2\beta_{\text{eff}} = 0$. In these measurements, both $B \to \eta'K^{0}_S$ and $B \to \eta'K^{0}_L$ final states are used. The $\eta'$ is reconstructed in its $\eta \pi^+\pi^-$ and $\rho^0\gamma$ decay modes. Combined, the analysis reconstruct more than 3000 signal events. Figure 6 shows the corresponding \textit{BABAR} event samples.

Using these events, \textit{BABAR} and Belle apply the same time-dependent analysis techniques as used in the $b \to c$ analysis above and determine

\begin{align}
\sin 2\beta_{\text{eff}} &= 0.58 \pm 0.10 \pm 0.03 \text{ (}\textit{BABAR}\text{),} \\
C &= -0.16 \pm 0.07 \pm 0.03 \text{ (}\textit{BABAR}\text{),} \\
\sin 2\beta_{\text{eff}} &= 0.64 \pm 0.10 \pm 0.04 \text{ (Belle),} \\
C &= 0.01 \pm 0.07 \pm 0.05 \text{ (Belle),}
\end{align}

both measuring $\sin 2\beta_{\text{eff}}$ to differ from 0 with greater than 5\sigma significance. No significant direct CP violation is observed in this mode. The $\Delta t$ distribution for $B \to \eta'K$ from the Belle analysis is shown in Fig. 7.
Figure 6: Event samples from the BABAR analysis of $B \rightarrow \eta'K_0^*$ (left and middle) and $B \rightarrow \eta'K_0^0$ (right). In each case, the points with error bars are the data, the dashed line illustrates the determined background level, and the solid line shows the total fit result.

Figure 7: $\Delta t$ for $B^0$ and $\bar{B}^0$ events (top) and raw asymmetry (bottom) distributions from the Belle analysis of $B \rightarrow \eta K$. The distributions for $B \rightarrow \eta'K_0^0$ have been flipped to account for the opposite $\eta'$ with respect to the $B \rightarrow \eta'K_0^0$ channel.
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LEPTON FLAVOUR VIOLATION

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A review of current experimental lepton flavour violation studies involving $\mu$ and $\tau$ decays is presented. Prospects for future experiments are also discussed.

1 Introduction

Historically, searches for lepton flavour violation (LFV) have had a tremendous impact on the field of particle physics, both when it was seen (in the neutrino sector\(^1\)) and unseen (in the charged lepton sector\(^2\)). Around the time that the muon was established to be a lepton, and not the Yukawa meson\(^3\), Hinks and Pontecorvo at Chalk River began searching for evidence of transitions of a charged lepton of one flavour into another via the emission of a photon\(^4\). Failure to observe this decay ruled out a model of the muon being an excited electron. In 1957 the lack of a flavour-changing neutral current in the lepton sector together with a “leptonic GIM” argument was used to establish the existence of two neutrino flavours\(^5,6,7\). Also in the 1950’s experimenters began searching for LFV in other processes. The first $\mu^- N \rightarrow e^- N$ neutrinoless muon nuclear capture experiments\(^8,9\) and searches for $\mu^- \rightarrow e^- e^+ e^-$ date from that era.

Since that time, the upper limits on branching fractions or conversion probabilities have fallen ten orders of magnitude from the percent level limits of the first studies to the “state-of-the-art” upper limit of $B(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11}$ from MEGA/LAMPF\(^10\). By the 1970’s searches for $\mu$-$e$ LFV in meson decays had also been performed and, in fact, new limits on LFV from kaons\(^11\) and b-mesons\(^12\) are being presented at this conference.

With the discovery of the $\tau$-lepton in 1974, a new LFV window opened and searches for a variety of lepton flavour violating decays were undertaken by the MARK-II, ARGUS and CLEO collaborations. These studies helped establish the SM description of the $\tau$ lepton: like its lighter cousin, is also is not an excited muon or electron and the existence of a third type of neutrino could be deduced. Clearly, the non-observation of LFV in the charged sector over the years has yielded critical input to the development of the standard model (SM).

In the neutral leptonic sector, however, a different story unfolds: by the beginning of the 21st century LFV has been established with the discovery of neutrino mixing\(^4\). In fact, we even have some knowledge of parts of the neutrino mixing matrix ($\theta_{12}$ and $\theta_{23}$) and plans for experiments to measure some of the unknown parts. For example, the T2K experiment is preparing to measure $\theta_{13}$, as discussed at this conference\(^13\) and long-term plans to measure CP violation in this sector are underway. The observation of LFV in this sector has generated a new field of exploration.

The SM, when extended to include finite mass differences from neutrino mixing, allows for $\mu - e, \tau - e$ and $\tau - \mu$ mixing at the one-loop level, but it is suppressed by a factor of $(m_\mu^2/m_\tau^2)^2$, which, for example, gives an expectation of $B(\tau \rightarrow \mu\gamma) \sim O(10^{-54})$, which is many orders of magnitude below any conceivable experimental sensitivity. As a consequence, this essentially non-existent SM background ensures that LFV in the charged sector is an unambiguous signature of new physics. This is very useful, in fact, since many SM extensions predict observable LFV in the charged sector. For example, if a see-saw mechanism is responsible for the small $\nu$ masses, it is natural to expect large LFV in the charged sector. In SO(10) inspired SUSY models there are regions of parameter space that allow for LFV above the existing limits\(^14\).
2 The LFV Programme Motivation

Needless to say, LFV predictions are very model dependent with specific models and parameters giving LFV process rates characteristic of those models. Generally speaking, we don't expect a single LFV process to provide sufficient information to nail down the underlying mechanism responsible for LFV, or even identify a particular model uniquely. Nonetheless, a strategy of combining results from many different, and varied measurements, is envisioned as the scientific programme that will move our understanding forward. For example, all the $\mu \to e$ LFV processes should be measured along with all $\tau$ LFV channels because many models have strong connections between the expected rates of the various channels. Consequently, the $\mu \to e\gamma$ measurements need to be augmented, for example, by $\tau \to \mu\gamma$ as well as $\tau \to e\gamma$ studies. Progress will be made by interpreting these results in the context of those of other measurements, including LFV decays of $K$ and $B$ mesons, neutrino oscillation measurements, $g-2$ and electric dipole moment measurements, and, naturally, searches for new particles at the energy frontier of LHC. Eventually precision measurements at a future ILC will also play an important role in this programme.

A very important step in the charged sector LFV component of these studies is to be realized later this year as the MEG experiment at Paul Scherrer Institut in Zurich takes first data.\cite{MEG} This project, which involves collaborators from Japan, Italy, Switzerland, Russia and the United States, aims at a $\mu \to e\gamma$ branching fraction sensitivity of $10^{-13}$ which is 100 times lower than the existing MEGA limit. The detectors are currently being built and installed with the first physics run slated for autumn 2007.

Since the MEGA results were published in 1999, new LFV results from charged lepton decays have come from studies of the third-generation lepton, which is the focus of the remainder of this paper. The LFV decays of the $\tau$ have been studied at $e^+e^-$ colliders where the $\tau$ leptons are most copiously pair-produced in an experimentally clean environment. The luminosities available at the $e^+e^-$ machines limit the number of $\tau$ leptons produced and thereby determine the sensitivity to $\tau$ LFV decays. Despite the unprecedented $10^{34} cm^{-1}s^{-1}$ luminosities attained the $e^+e^-$ B-factories, these luminosity-limited sensitivities are still typically 1000 times weaker than the MEGA $\mu \to e\gamma$ upper limit. Given that reality, one may wonder if studies of $\tau$ decays bring anything new to the table. The answer is yes: many SM extensions predict very sizable enhancements of the $\tau$ LFV decays relative to the $\mu$ LFV processes. These usually originate from strong dependencies of the couplings on the masses of the leptons involved, so that a LFV process involving a $\tau$-$\mu$ transition, which involves a $1777$ MeV mass lepton converting to a $106$ MeV lepton, may be greatly enhanced relative to a process involving a $\mu - e$ transition which converts a $106$ MeV lepton to a $0.5$ MeV lepton. Consequently, one encounters theories with predictions of $B(\tau \to \mu\gamma)$, for example, that have parameters being constrained by the existing experimental bounds.

Equally important is the model dependence of the expected rates for different $\tau$ LFV decays and the relations between these rates. For example, in a supersymmetric seesaw model there is an expectation that the specific relative rates of $B(\tau \to \mu\gamma) : B(\tau \to \mu\mu\mu) : B(\tau \to \mu\eta)$ are dependent on the model parameters\cite{appellof,appellof2}. Because we do not know what lies beyond the SM, it is unknown which LFV decay mode will be first discovered and therefore it is critical that all LFV modes be probed. Limits, or (preferably) a discovery and measurement, will better constrain theories in a manner that complements well the potential discoveries at the LHC.

3 The $\tau$ LFV Experiments

In recent years, only the Babar and Belle experiments have been providing new results on LFV in $\tau$ decays. These operate at two different $e^+e^-$ "B-factories" both tuned to a centre-of-mass
energy of 10.58 GeV: at the $\Upsilon(4S)$ resonance, and just above the threshold for producing open beauty. Babar operates at the PEP-II B-factory at SLAC, running the $e^-$ beam at 9 GeV and $e^+$ beam at 3.1 GeV. By 1 January 2007 Babar had written more that 400 fb$^{-1}$ of data to tape. Belle, operating at the KEKB B-factory in Japan, which collides an 8 GeV $e^-$ beam with a 3.5 GeV $e^+$ beam, has recorded an integrated luminosity of approximately 700 fb$^{-1}$. The Babar$^{18}$ and Belle$^{19}$ detectors are remarkably similar with the major difference being in the technology used to identify charged particles: Belle uses a threshold Cherenkov detector together with time-of-flight and tracker dE/dx, whereas Babar mainly relies on a ring-imaging Cerenkov detector augmented by the dE/dx in the trackers. With a total of 1.1 ab$^{-1}$ and the $e^+e^-\rightarrow \tau^+\tau^-$ cross-section of 0.919 nb$^{20}$, the world sample of $\tau$-leptons produced at the $e^+e^-$ colliders now exceeds $10^3$ which allows for experimental probing of LFV processes at the $O(10^{-7})$ to $O(10^{-8})$ levels.

The general analysis approach is to select $\tau$-pair events with the appropriate charged-particle topology, removing non-$\tau$ events with an impact as minimal as possible on the signal efficiency. This is accomplished by dividing the candidate event into hemispheres in the centre-of-mass in which each hemisphere contains either the $\tau^+$ or $\tau^-$ decay products. Each hemisphere is then considered a possible candidate for the LFV decay under consideration. Unlike SM $\tau$-decays which have at least one neutrino, the LFV decay products have a combined energy in the centre-of-mass equal to the energy of the $\tau$, approximately equal to the beam energy in the centre-of-mass, $\sqrt{s}/2$, and a mass equal to that of the $\tau$. A two dimensional signal box in $E_{\ell\chi}$ vs. $M_{\ell\chi}$ is therefore used to separate the signal from the SM $\tau$-decay backgrounds. The distribution of $\Delta_{E} = E_{\ell\chi} - \sqrt{s}/2$ for the signal peaks near zero and typically has a standard deviation of around 50 MeV. When a beam energy constrained mass is employed and the $\gamma$ constrained to come from the same primary vertex as the charged particles in the event, then a resolution on $M_{\ell\chi}$ of 9 MeV is achieved.

The analyses are optimized to give the best “expected upper limit” using Monte Carlo simulations of the signal and backgrounds. The Monte Carlo simulation of the signal is used to determine the signal efficiency ($\epsilon$), which typically lies between 2% and 10%, depending on the channel under study. A generic analysis of a $\tau$ LFV decay processes roughly the following efficiency components: trigger (90%), acceptance/reconstruction (70%), charged-particle hemisphere topology (1-vs-1 or 1-vs-3: 70%), particle identification (50%), other non-signal box requirements (50%), $E_{\ell\chi}$ vs. $M_{\ell\chi}$ signal box requirements (50%). The cumulative efficiencies, starting with the trigger and acceptance/reconstruction efficiency of 63%, decreases to 44% once topology requirements are imposed, to 22% with particle ID requirements, to 11% as other non-signal box requirements are imposed, and finally to $\sim 5\%$ once it is enforced that the events fall within signal-box.

The expected number of background events ($N_{\text{bkd}}$) are normally estimated using the distribution shapes from the Monte Carlo simulation with background normalization obtained from the data in the regions outside the signal box in the $(E_{\ell\chi}-M_{\ell\chi})$ plane. These analyses are “blind” in the sense that the physics analysts have no knowledge of the data in the signal region as the optimization and systematic studies are undertaken. Once these steps are complete, the data in the signal region is “unblinded” and the analyst sees the number of events observed in the signal box ($N_{\text{obs}}$) and either learns of a discovery, or - as has been the case to date - sets an upper limit on the process.

$N_{\text{obs}}$ together with $N_{\text{bkd}}$ then gives the number of signal events ($N_{\text{sig}}$): if $N_{\text{obs}}-N_{\text{bkd}}$ is consistent with zero, an upper limit on $N_{\text{sig}}$ is established. Schematically, the 90%CL branching ratio upper limit is then obtained from:

$$B_{90}^{UL} = \frac{N_{90}^{UL}}{2N_{\tau\tau}\epsilon} = \frac{N_{90}^{UL}}{2\mathcal{L}_{\tau\tau}\epsilon}$$  \hspace{1cm} (1)
Table 1: Summary of 90%CL Upper Limits on Selected LFV \( \tau \) decays. An asterix\((*)\) indicates a preliminary result. \( h \) and \( h' \) denotes a charged pion or kaon. Banerjee's combination of a subset of these channels is also included - note that for the combined results only, the limits are in units of \( 10^{-8} \).

<table>
<thead>
<tr>
<th>Channel</th>
<th>Belle 21,24,25,26</th>
<th>Babar 22,27,28,29,30</th>
<th>Babar+Bel 23</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \mu\gamma )</td>
<td>0.5*</td>
<td>0.7</td>
<td>1.6</td>
</tr>
<tr>
<td>( \mu\pi^0 )</td>
<td>1.2*</td>
<td>1.1</td>
<td>5.8</td>
</tr>
<tr>
<td>( \mu\eta )</td>
<td>0.7*</td>
<td>1.5</td>
<td>5.1</td>
</tr>
<tr>
<td>( \mu\eta' )</td>
<td>1.3*</td>
<td>1.4</td>
<td>5.3</td>
</tr>
<tr>
<td>( e\gamma )</td>
<td>1.2*</td>
<td>1.1</td>
<td>9.4</td>
</tr>
<tr>
<td>( e\pi^0 )</td>
<td>0.8*</td>
<td>1.3</td>
<td>4.4</td>
</tr>
<tr>
<td>( e\eta )</td>
<td>0.9*</td>
<td>1.6</td>
<td>4.5</td>
</tr>
<tr>
<td>( e\eta' )</td>
<td>1.6*</td>
<td>2.4</td>
<td>9.0</td>
</tr>
<tr>
<td>( lll )</td>
<td>2→4</td>
<td>1→3</td>
<td>92</td>
</tr>
<tr>
<td>( llh )</td>
<td>2→16</td>
<td>1→5</td>
<td>221</td>
</tr>
</tbody>
</table>

where \( N_{\tau\tau} = \mathcal{L} \sigma_{\tau\tau} \) is the number of \( \tau \)-pairs produced in \( e^+ e^- \) collisions; \( \mathcal{L} \) is the integrated luminosity and \( \sigma_{\tau\tau} \) is the \( \tau \)-pair production cross section. In practice, if \( N_{\text{bkgd}} \) is sufficiently large (more than approximately two or three), \( N_{\text{sig}} \) and \( N_{\text{bkgd}} \) are determined from a fit to the \( M_{lX} \) distribution after removing events outside the \( E_{lX} \) signal region.

4 The \( \tau \) LFV Results

Experimentally, LFV decays can be conveniently classified as \( \tau \to \ell\gamma \), \( \tau \to \ell_1\ell_2\ell_3 \) and \( \tau \to \ell h \) where \( \ell \) is either an electron or muon and \( h \) represents a hadronic system (e.g., \( \pi^0 \), \( \eta \), \( \eta' \), \( K^0_S \), etc.)

The most recent \( \tau \to \mu\gamma \) and \( \tau \to e\gamma \) results were reported by Belle at ICHEP last summer\(^{21}\). The \( \tau \to \mu\gamma \) \((\tau \to e\gamma)\) analyses have a 5.1\% (3\%) signal efficiency within a 2\( \sigma \) elliptical signal region in the \( M_{\ell\gamma} - E_{\ell\gamma} \) plane. After performing a 2D unbinned extended maximum likelihood fit for the number of signal events and background events, Belle finds 3.9 (0.14) signal events and 13.9 (5.14) background events in the \( \tau \to \mu\gamma \) \((\tau \to e\gamma)\) analysis. Therefore the 90\%CL upper limits on the number of signal events for \( \tau \to \mu\gamma \) \((\tau \to e\gamma)\) is 2.0 (3.34) events. These yield upper limits of \( B(\tau \to \mu\gamma) < 4.5 \times 10^{-8} \) and \( B(\tau \to e\gamma) < 1.2 \times 10^{-7} \).

Babar has very recently published new results on LFV decays involving the \( \pi^0 \), \( \eta \) and \( \eta' \) pseudoscalars; \( \tau \to \ell\pi^0 \), \( \tau \to \ell\eta \) and \( \tau \to \ell\eta' \) where \( \ell \) is separately identified as either an electron or muon\(^{22}\). In these analyses both of the \( \eta \to \gamma\gamma \) and \( \eta \to 3\pi \) \( \eta \) decay modes are used for the \( \tau \to \ell\eta \) analyses. In the \( \tau \to \ell\eta' \) analyses, the \( \eta' \to 2\pi\eta \) and \( \eta' \to 2\pi\gamma \) modes were included. The expected background per channel is between 0.1 and 0.3 events. Summing over all ten modes, the total expected background within the signal regions is 3.1 events whereas 2 events in total were observed. The results of these searches and a selection of searches for other key LFV decays in data from Babar and Belle are summarized in Table 1. At the Tau2006 conference in Pisa, Swagata Banerjee presented a frequentist combination of these measurements\(^{23}\) which are also included in the table.

5 Future Prospects for \( \tau \) LFV Searches

Belle has funding secured that will enable it to collect 1~ab\(^{-1}\) of data and Babar, which will stop data taking at the end of September 2008, is expected to collect 940~fb\(^{-1}\) of data. The
estimated physics reach of these data based on projections from existing analyses depends on how the background is treated. We express the experimental reach in terms of the ‘expected 90%CL upper limit’ and, for brevity’s sake, refer to this as the ‘sensitivity’. In the absence of signal, for large $N_{\text{bkd}}$, $N_{90}^{UL} \approx 1.64\sqrt{N_{\text{bkd}}}$ whereas for small $N_{\text{bkd}}$ a value for $N_{90}^{UL}$ is obtained using the method described in \cite{31}. So, for $N_{\text{bkd}} \sim 0$, $N_{90}^{UL} \sim 2.4$. Reducing the background below a handful of events doesn’t greatly improve the expected limit if significant efficiency is lost in the process, which is why it is common to see experiments reporting the expected backgrounds to be small (i.e. a few events), but rarely below 0.1 of an event.

A “worst-case scenario” is obtained if identical analyses to those published by Babar and Belle are repeated, as is, on the increased data sample: in that case the expectations then simply scale as $\sim \sqrt{N_{\text{bkd}}}/\mathcal{L}$, which for large $N_{\text{bkd}}$ scales as $1/\sqrt{\mathcal{L}}$. A “best case” scenario would take the current expected limit and scale linearly with the luminosity. This is equivalent to a statement that analyses can be developed maintaining the same efficiency and backgrounds as the current analyses.

For $\tau \to \ell\gamma$, there is an “irreducible background” from $\tau \to \ell\nu\nu + \gamma$ (ISR) in which the photon from initial state radiation can be combined with a lepton to form a candidate that accidentally overlaps with the signal region in the $E_{\text{T}} - M_{\text{T}}$ plane. It is “irreducible” in the sense that it arises from $e^+e^- \to \tau^+\tau^-$ process with a well measured $\mu$ and $\gamma$ in one of the $\tau$ hemispheres. In the existing Babar analyses, these events account for approximately one fifth of the total background. Scaling with this irreducible background only, one has an expected upper limit for $B(\tau \to \ell\gamma)$ of between 1 and $2 \times 10^{-5}$ from a complete combined Babar and Belle data set. That Banerjee’s combined limit on $B(\tau \to \mu\gamma)$ is already at this level is a consequence of a downward statistical fluctuation in the expected background.

The situation for the other LFV decays, $\tau \to \ell_1\ell_2\ell_3$ and $\tau \to \ell h$, is even more promising, since these modes do not suffer from the aforementioned backgrounds from ISR. In this case, one can project sensitivities assuming $N_{\text{bkd}}$ comparable to backgrounds in existing analyses for approximately the same efficiencies. These yield expected limits at the $10^{-8}$ level with the complete Babar+Belle data set.

It should also be mentioned that there are proposals \cite{32} for Super B flavour factories which will generate up to a 100 fold increase in the size of the $\tau$ sample compared to those expected from the existing B-factories. If such a facility is built, one will be probing LFV decays of the $\tau$ at the $\mathcal{O}(10^{-9} - 10^{-10})$ level.

6 Summary

Lepton flavour violation in the decays of charged leptons is an extremely clean means of searching for evidence of physics beyond the SM. Until now, there has been no evidence of LFV in either the $\mu \to e\gamma$ decay nor in any of a number of $\tau \to \mu X$ or $\tau \to eX$ decays, where $X$ is a photon, meson or set of mesons. Recent results from Babar and Belle have pushed the upper limits on the branching fractions of these decays of the $\tau$ into the $10^{-8}$ zone and the parameter space in some beyond-the-SM theories is being constrained. As there is still much data yet to be collected and analysed by the B-factory experiments, we look forward to continuing progress on this front over the next two years.

With physics data from MEG to come online this autumn, we also look forward to progress on the $\mu \to e\gamma$ front, ultimately, with sensitivities at the $\mathcal{O}(10^{-13})$ level being reached. In the more distant future, if Super B flavour factories move forward, sensitivities to some $\tau$ LFV modes will approach the $\mathcal{O}(10^{-10})$ level. Probing LFV at these tiny levels will provide critical information about the nature of the beyond-the-SM theories in a manner that is complementary to that from direct production of new particles at the LHC.
Acknowledgments

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Semileptonic $B$ decays at the $B$ factories

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I present in this review recent results on semileptonic $B$ decays at the $B$ factories.

1 Introduction

Semileptonic $B$ decays to charm and charmless mesons are the primary tool where to measure the CKM matrix elements $|V_{cb}|$ and $|V_{ub}|$ because of their simple theoretical description at the parton level. Their relatively large decay rates are proportional to $|V_{cb}|^2$ and $|V_{ub}|^2$ and depend on the quark masses $m_b$, $m_c$. In recent years, tremendous progress has been made in the description of semileptonic and radiative decays in the framework of Heavy Quark Expansions (HQEs); calculations for their width and moments of inclusive observables with restrictions on the phase space are now available in different schemes through order $1/m_b^2$.

2 HQE Parameters Determination

Making use of HQEs it is possible to perform a combined fit to measured moments for which correlation matrices are published. These include moments of the hadronic mass distribution $\langle M_X^2 \rangle$ and moments of the lepton energy spectrum $\langle E_l^m \rangle$ in inclusive $B \to X_c \ell \bar{\nu}$ as well as moments of the photon energy spectrum $\langle E_\gamma^m \rangle$ in inclusive $B \to X_\gamma \gamma$ decays for different minimum lepton and photon energies. The important issue in this approach is the ability to separate non-perturbative corrections (expressed as a series in powers of $1/m_b$), and perturbative corrections (expressed in powers of $\alpha_S$). There are various different methods to handle the energy scale $\mu$ used to separate long-distance from short-distance physics.

There are two schemes that will be considered in this paper: Kinetic$^2$ and $1S^1$ scheme.
2.1 Experimental Input

Belle has measured the partial branching fractions \( B(B \rightarrow X_c\ell\bar{\nu})_{E_l > E_{\text{min}}} \) and the first, second, third, and fourth moments of the truncated electron energy spectrum in \( B \rightarrow X_c\ell\bar{\nu} \), \( \langle E_l \rangle_{E_l > E_{\text{min}}} \), \( \langle (E_l - \langle E_l \rangle)^2 \rangle_{E_l > E_{\text{min}}} \), \( \langle (E_\ell - \langle E_\ell \rangle)^3 \rangle_{E_l > E_{\text{min}}} \), and \( \langle (E_\ell - \langle E_\ell \rangle)^4 \rangle_{E_l > E_{\text{min}}} \), for nine different lepton energy thresholds \( E_{\text{min}} = 0.4, 0.6, 0.8, 1.0, 1.2, 1.4, 1.6, 1.8 \) and 2.0 GeV. Belle has also measured the first, second central and second non-central moments of the hadron invariant mass squared \( (M_X^2) \) spectrum in \( B \rightarrow X_c\ell\bar{\nu} \), \( \langle M_X^2 \rangle_{E_l > E_{\text{min}}} \), \( \langle (M_X^2 - \langle M_X^2 \rangle)^2 \rangle_{E_l > E_{\text{min}}} \), and \( \langle M_X^2 \rangle_{E_l > E_{\text{min}}} \) for seven different lepton energy thresholds \( E_{\text{min}} = 0.7, 0.9, 1.1, 1.3, 1.5, 1.7 \) and 1.9 GeV. \(^4\) For \( B \rightarrow X_s\gamma \), Belle has measured the first and second moments of the truncated photon energy spectrum, \( \langle E_\gamma \rangle_{E_l > E_{\text{min}}} \) and \( \langle (E_\gamma - \langle E_\gamma \rangle)^2 \rangle_{E_l > E_{\text{min}}} \), for six minimum photon energies \( E_\gamma = 1.8, 1.9, 2.0, 2.1, 2.2 \) and 2.3 GeV. \(^5\) Out of the 71 measurement of inclusive spectra, the ones that have no corresponding theoretical prediction were not used to determine HQE parameters, along with the ones with higher cutoff energies (i.e. electron energy and hadron mass moment with \( E_{\text{min}} > 1.5 \) GeV and photon energy moment with \( E_{\text{min}} > 2 \) GeV). Points where correlations with neighboring points are too high were also excluded.

2.2 Kinetic Scheme

Inclusive spectral moments of \( B \rightarrow X_c\ell\bar{\nu} \) have been derived in the Kinetic scheme up to \( \mathcal{O}(1/m_b^3) \). Theoretical expressions of these moments depend on the b- and c-quark masses \( m_b(\mu) \) and \( m_c(\mu) \), the non-perturbative parameters \( \mu_2^2(\mu) \) and \( \mu_3^2(\mu) \) \( \mathcal{O}(1/m_b^2) \), \( \psi_s(\mu) \) and \( \psi_s(\mu) \) \( \mathcal{O}(1/m_b^3) \), and \( \alpha_s \). A \( \chi^2 \) function can be written, that combines the measured moments and the corresponding kinetic scheme predictions that depend on these free parameters. A fit that minimizes the \( \chi^2 \) function and that takes into account both experimental and theoretical uncertainties was performed, where the CKM matrix element \( |V_{cb}| \) was treated as an eighth free parameter. The fit results are summarized in the following\(^6\):

\[
|V_{cb}| = (41.93 \pm 0.65_{\text{fit}} \pm 0.07_{\alpha_s} \pm 0.63_{\text{th}}) \times 10^{-3},
\]
\[
m_B = 4.564 \pm 0.076_{\text{fit}} \pm 0.003_{\alpha_s} \text{GeV},
\]
\[
m_c = 1.105 \pm 0.116_{\text{fit}} \pm 0.005_{\alpha_s} \text{GeV},
\]
\[
B(B \rightarrow X_c\ell\bar{\nu}) = 10.590 \pm 0.164_{\text{fit}} \pm 0.006_{\alpha_s} %,
\]

where first error is due to all uncertainties taken into account in the fit (experimental error in the moment measurements, non-perturbative corrections and bias correction to the moments, uncertainty in \( \tau_B \)), the second error comes from the uncertainty in \( \alpha_s \), and the last error in the expression of \( |V_{cb}| \) is a 1.5% uncertainty due to the limited accuracy of the theoretical expression for the semileptonic width. The fit results for \( m_b, \mu_2^2 \) and \( \rho_3^3 \) are shown in Figure 1.

2.3 1S scheme

The inclusive spectral moments of \( B \rightarrow X_c\ell\bar{\nu} \) have also been derived in the 1S scheme up to \( \mathcal{O}(\alpha_s/m_b^3) \). A total of seven parameters were determined: \( |V_{cb}|, \Lambda, \lambda_1, \tau_1, \tau_2, \tau_3 \) and \( \rho_1 \). The parameter \( \Lambda \) is the difference between the b-quark mass and the reference value about which it is expanded, i.e., \( \Lambda = m_{(1S)}/2 - m_b^{1S} \), and results will be quoted in terms of \( m_b^{1S} \). A fit to a \( \chi^2 \) function that combines the measured moments and the corresponding 1S scheme predictions and that takes into account both experimental and theoretical uncertainties was performed, giving the results\(^6\):

\[
|V_{cb}| = (41.49 \pm 0.52_{\text{fit}} \pm 0.20_{\tau_B}) \times 10^{-3},
\]
\[
m_b^{1S} = (4.729 \pm 0.048) \text{ GeV},
\]
\[
\lambda_1 = (0.30 \pm 0.04) \text{ GeV}^2,
\]
Figure 1: Fit results for $m_b$ and $\mu^2$ (left) and $m_b$ and $\rho_B^2$ in the Kinetic Scheme. The fit is done using lepton energy moments only (blue), hadron mass moments only (yellow), photon energy moments (green) only and all moments combined (red). The ellipses correspond to $\Delta \chi^2 = 1$.

Figure 2: Fit results for $m_b^{1S}$ and $\lambda_1$ (left) and $m_b^{1S}$ and $|V_{cb}|$ (right) in the 1S Scheme. The fit to the $B \rightarrow X_c \ell \bar{\nu}$ data only ($B \rightarrow X_c \ell \bar{\nu}$ and $B \rightarrow X_{c\gamma}$ data combined) is shown by a dashed blue line (solid red line). The regions correspond to $\Delta \chi^2 = 1$.

where the first error is the uncertainty from the fit including experimental and theory errors, and the second error (on $|V_{cb}|$ only) is due to the uncertainty on the average $B$ lifetime. The fit results for $m_b^{1S}$, $\lambda_1$ and $|V_{cb}|$ are shown in Figure 2.

3 Relative Branching Fractions for $B \rightarrow D, D^*, D^{**} \ell \nu$ Decays

The determination of the exclusive branching fractions of $B \rightarrow X_c \ell \bar{\nu}$ decays is an essential part of the program of studies of semileptonic B meson decays at the $B$-factories to understand the dynamics of semileptonic decays of the $b$ quark and determine the relevant CKM matrix elements. The mass of the hadronic system $X_c$, recoiling against the fermionic pair in the decay, is a crucial parameter both in the extraction of $|V_{cb}|$, in exclusive semileptonic decays, and in isolating $B \rightarrow X_{c\ell} \ell \bar{\nu}$ decays for determining $|V_{ub}|$. It is also important in the measurement of the heavy quark masses and other non-perturbative OPE parameters from the distribution of spectral moments in semileptonic decays.

$B\bar{B}$ has presented a novel technique $^7$ to extract the exclusive relative branching fractions for $B \rightarrow D \ell \bar{\nu}$, $B \rightarrow D^* \ell \bar{\nu}$ and $B \rightarrow D^{**} \ell \bar{\nu}$ with $\ell = e, \mu$ in an inclusive sample of $B \rightarrow D X \ell \bar{\nu}$ events, where X can be either nothing or any particle(s) from a semileptonic $B$ decay into a higher mass charm state, or a non-resonant state. $D^{**}$ is a hadronic final state containing a charm meson, with a total mass above the $D^*$ mass, therefore including both $D^{**}$ mesons and non-resonant $D^{(*)} + n \pi$ states.

The analysis is based on 339.4 $fb^{-1}$ of data collected at the $\Upsilon$(4S) resonance with the $B\bar{B}$
detector at the PEP-II asymmetric-energy $e^+e^-$ storage rings. Signal $B$ mesons are selected by fully reconstructing the accompanying $B$ meson ($B_{tag}$), which allows to constrain the kinematics, to reduce the combinatorial background and to determine the charge of the signal $B$.

A set of three largely uncorrelated variables to discriminate between the different semileptonic decay modes in the reconstructed $B \rightarrow DX\ell\bar{\nu}$ sample was chosen. These are: the lepton momentum in the center-of-mass frame, $p_\ell$; the missing mass squared reconstructed with respect to the $D\ell$ system, which corresponds to the mass of the $X\ell\bar{\nu}$ system, $m_{miss,D}^2 = (p_T - p_{D_{tag}} - p_D - p_\ell)^2$; the number of reconstructed charged tracks in addition to those used for reconstructing the $D\ell$ system and the $B_{tag}$, $N_{trks}$. In order to reduce the sensitivity to the modeling of the decays to the different charm states, the shapes of these variables were extracted from the data itself, using samples highly enriched in the relevant decay modes. The relative $D$, $D^*$, $D^{**}$ contributions were then determined by a multi-parameter fit to the inclusive sample. The result of the fit to the missing mass distribution for $B^0 \rightarrow DX\ell\bar{\nu}$ is shown in Figure 3. The relative branching fractions for $B^- \rightarrow D^0, D^{*0}, D^{**0}\ell\bar{\nu}$ and $B^0 \rightarrow D^+, D^{*+}, D^{**+}\ell\bar{\nu}$, with statistical and systematic uncertainties are given in Table 1.

4 Simultaneous Determination of the $B^0 \rightarrow D^{*\pm}\ell^+\nu_{\ell}$ Form Factors and of $|V_{cb}|$

In the Standard Model, the rate of the $B^0 \rightarrow D^{*\pm}\ell^+\nu_{\ell}$ decay is proportional to $|V_{cb}|^2$. The decay is also influenced by strong interactions, parameterized by two axial form factors $A_1$ and $A_2$, and one vector form factor $V$, each of which depends on the momentum transfer squared $q^2$ of the $B$ meson to the $D^*$ meson. The form of this dependence is not known a priori. In the framework of heavy-quark effective field theory (HQET)\textsuperscript{8,9}, these three form factors are related to each other through heavy quark symmetry (HQS), but HQET leaves three free parameters, which must be determined by experiment. The form factors are usually characterized in terms of their ratio parameters $R_1$ and $R_2$, and a slope parameter $\rho^2$.

The extraction of $|V_{cb}|$ relies on the measurement of differential decay rates. HQS predicts the

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**Table 1: Fitted ratios of branching fractions with statistical and systematic uncertainties**

<table>
<thead>
<tr>
<th>Ratio</th>
<th>$B^-$ (%)</th>
<th>$B^+$ (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Gamma(B \rightarrow D\ell\bar{\nu})$</td>
<td>$22.7 \pm 1.4 \pm 1.6$</td>
<td>$21.5 \pm 1.6 \pm 1.3$</td>
</tr>
<tr>
<td>$\Gamma(B \rightarrow D^{*\ell}\bar{\nu})$</td>
<td>$58.2 \pm 1.8 \pm 3.0$</td>
<td>$53.7 \pm 3.1 \pm 3.6$</td>
</tr>
<tr>
<td>$\Gamma(B \rightarrow D^{**\ell}\bar{\nu})$</td>
<td>$19.1 \pm 1.3 \pm 1.9$</td>
<td>$24.8 \pm 3.2 \pm 3.0$</td>
</tr>
</tbody>
</table>
normalization of the decay rate at the maximum $q^2$, and $|V_{cb}|$ is determined from an extrapolation of the form factors to this value. The precise determination requires corrections to the HQS prediction for the normalization as well as a measurement of the variation of the form factors near the maximum $q^2$, where the decay rate goes to zero as the phase space vanishes.

79 fb$^{-1}$ of $B \overline{B}$ events collected with the BaBar detector were analyzed$^{10}$; in place of a linear parameterization of the $q^2$ dependence of the form factors, a higher-order polynomial was used. The Lorentz structure of the $B \rightarrow D^*\ell\nu$ decay amplitude can be expressed in terms of three helicity amplitudes which correspond to the three polarization states of the $D^*$. The full differential decay rate can be expressed in terms of four kinematic variables: the Lorentz-invariant $w$ (the dot product of the $B$ and $D^*$ four-velocities) and three angles $\theta_L$ (the angle between the direction of the lepton in the virtual $W$ rest frame, and the direction of the virtual $W$ in the $B$ rest frame), $\theta_W$ (the angle between the direction of the $D$ in the $D^*$ rest frame) and $\chi$ (the dihedral angle between the plane formed by the $D^* - D$ and the plane formed by the $W - \ell$ system).

$|V_{cb}|$ and the form factors were extracted by performing a $\chi^2$ fit of three one-dimensional distributions, as it was found that the $\chi$ distribution is insensitive to the form factor parameters. The three other kinematic variables were divided into ten bins, and a least-square fit of the projected one-dimensional distributions was performed, correcting for background contributions. The result of the fit is shown in Figure 4.

Figure 4: Comparison of the measured distributions (data points) a) $w$, b) $\cos \theta_L$, c) $\cos \theta_W$, and d) $\chi$, with the result of the fit, shown as the sum of the fitted signal yield and the estimated background distributions. The statistical uncertainties of the data are too small to be visible.

The results of the extraction of $|V_{cb}|$ and the form factors were then combined with the results of an earlier BaBar measurement$^{11}$, based on an unbinned maximum-likelihood fit to the four-dimensional decay distribution; this fit is sensitive to the interference of the three helicity amplitudes and thus results in significant smaller uncertainties on the form-factor parameters.
Table 2: Partial branching fractions in three bins of \(q^2\). These are summed to give the full branching fraction quoted in the Sum column. Errors are statistical and systematic.

<table>
<thead>
<tr>
<th>Mode</th>
<th>(\Delta B (10^{-4}))</th>
<th>(B (10^{-4}))</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>(0 &lt; q^2 &lt; 8 \text{ (GeV}/c^2))</td>
<td>(8 &lt; q^2 &lt; 16 \text{ (GeV}/c^2))</td>
</tr>
<tr>
<td>(B \rightarrow \pi^+ \ell \nu)</td>
<td>0.50 ± 0.14 ± 0.02</td>
<td>0.68 ± 0.18 ± 0.03</td>
</tr>
<tr>
<td>(B \rightarrow \pi^0 \ell \nu)</td>
<td>0.28 ± 0.09 ± 0.02</td>
<td>0.22 ± 0.08 ± 0.04</td>
</tr>
</tbody>
</table>

Taking into account the correlation between the two measurements, \(\Lambda\bar{B}\Lambda\bar{B}\) obtains:

\[
\mathcal{F}(1) |V_{ub}| = (34.7 \pm 0.3 \pm 1.1) \times 10^{-3}
\]

\[
\rho^2 = 1.179 \pm 0.048 \pm 0.028
\]

\[
R_1(1) = 1.417 \pm 0.061 \pm 0.044
\]

\[
R_2(1) = 0.836 \pm 0.037 \pm 0.022
\]

5 Branching Fractions for Exclusive \(B \rightarrow X_u \ell \bar{\nu}\) Decays

Exclusive charmless semileptonic decays can be used to extract the CKM matrix element \(|V_{ub}|\).

It is the aim of an ongoing programme of measurements at the \(B\) factories to improve the determination of \(|V_{ub}|\) with a precision better than 5%, for comparison with the inclusive results, which have somewhat different experimental and theoretical uncertainties.

The Belle collaboration has recently presented studies of the exclusive semileptonic decays \(B \rightarrow \pi^+ \ell \nu\), \(B \rightarrow \pi^0 \ell \nu\), \(B \rightarrow \rho^+ \ell \nu\), \(B \rightarrow \rho^0 \ell \nu\) and \(B \rightarrow \omega \ell \nu\) using a full reconstruction tagging technique to identify candidate \(B\) mesons\(^1\). The measurement presented here is based on a data sample that contains \(535 \times 10^6 \bar{B}B\) pairs, collected with the Belle detector at the KEKB asymmetric-energy \(e^+e^-\) (3.5 on 8 GeV) collider. One of the two \(B\) mesons from the pair is fully reconstructed in hadronic modes, allowing to constrain the kinematics, to reduce the combinatorial background and to determine the charge of the signal \(B\). Signal events are identified by examining the missing mass squared \((m_{\text{miss}}^2)\) distributions. If the tagging \(B\) is correctly reconstructed and the correct lepton and hadron candidate have been identified on the recoil side, then (ideally) all missing 4-momentum is due to the remaining unidentified neutrino. The square of the missing 4-momentum for signal events should therefore be close to zero, and applying this requirement provides a very strong discrimination between signal and background.

Fitting the \(m_{\text{miss}}^2\) distributions, it is possible to identify signal and background components. The fitted number of signal was: \(48 \pm 8\) for \(B \rightarrow \pi^+ \ell \nu\), \(35 \pm 7\) for \(B \rightarrow \pi^0 \ell \nu\), \(41 \pm 9\) for \(B \rightarrow \rho^+ \ell \nu\), \(61 \pm 9\) for \(B \rightarrow \rho^0 \ell \nu\) and \(27 \pm 9\) for \(B \rightarrow \omega \ell \nu\). Partial branching fractions were extracted in bins of \(q^2\), the invariant mass squared of the lepton-neutrino system, in order to minimize the systematic error which arises from the lack of precise knowledge of the shape of the form factors. Three bins in \(q^2\) were chosen, 0 to 8 GeV\(^2\)/c\(^2\), 8 to 16 GeV\(^2\)/c\(^2\), and greater than 16 GeV\(^2\)/c\(^2\).

The results are summarized in table 2.

6 Extraction of \(|V_{ub}|\) with Reduced Model Dependence

Theoretically, the most precise predictions can be made for the total \(B \rightarrow X_u \ell \bar{\nu}\) decay rate. Accounting for restriction in phase space is difficult because decay spectra close to the kinematic limit are susceptible to non-perturbative strong-interaction effects. Theoretical tools for the calculations of the partial inclusive decay rates are QCD factorization and local operator product
expansions (OPE). These calculations separate non-perturbative from perturbative quantities, and use expansions in inverse powers of the b-quark mass, and in powers of \( \alpha_S \). Non-perturbative bound-state effects are accounted for by a shape function describing the Fermi-motion of the b quark inside the B meson. These shape functions cannot be calculated. Different shapes have been proposed, and parameters defining these shapes have to be extracted from data. This introduces significant additional hadronic uncertainties. At leading order, these shape functions are assumed to be universal functions for \( b \to q \) transitions, where \( q \) is a light quark, either \( s \) or \( u \). Usually, the non-perturbative parameters of these shape functions, the b-quark mass \( m_b \) and its mean kinetic energy squared, \( \mu^2(\mu) \), are extracted from the inclusive photon spectrum in \( B \to X_s \gamma \) decays. These parameters depend on the choice of the renormalization scale, \( \mu \). It was suggested that \( |V_{ub}| \) can be extracted with smaller theoretical uncertainties by combining integrals over the lepton-endpoint spectrum in \( B \to X_u \ell \bar{\nu} \) decays with weighted integrals over the photon-endpoint spectrum from \( B \to X_s \gamma \) decays. The underlying assumption is that the QCD interactions affecting these two processes are the same and thus will cancel to first order in the appropriate ratio of weighted rates. The advantage of this approach is that it reduces the sensitivity to the choice of the shape-function parameterization and thus avoids uncertainties that are difficult to quantify.

**BABAR** has presented the extraction of \( |V_{ub}| \)\(^{16} \) based on different prescriptions proposed by Leibovich, Low, and Rothstein (LLR)\(^{13} \), by Neubert\(^{14} \), and more recent calculations by Lange, Neubert, and Paz (LNP)\(^{15} \).

The experimental inputs for this analysis are the published **BABAR** measurements of the inclusive electron energy spectrum in \( B \to X_u \ell \bar{\nu} \) decays\(^{17} \), and of the inclusive photon-energy spectrum in \( B \to X_s \gamma \) decays\(^{18} \). These measurements are based on a data sample corresponding to a total integrated luminosity of about 80 fb\(^{-1} \). The measured spectra were integrated above an energy \( E_0 \), measured in the B-meson rest frame.

A comparison of the extraction of \( |V_{ub}| \) for the three methods considered, with its dependence on the lepton energy cut-off, is shown in Figure 5. The LNP method is more accurate, in the sense that it includes two-loop calculations, and this explains the error inflation that we notice at large lepton energy cut-off. The three methods agree well at low lepton energy cut-off and give results with a precision comparable to Shape Function based analyses. At present, for \( E_0 = 2.0 \) GeV, the experimental errors on both the \( B \to X_u \ell \bar{\nu} \) branching fraction and the integral over \( B \to X_s \gamma \) are about 12%; this means that there are opportunities to improve the accuracy of \( |V_{ub}| \) significantly by reducing the experimental errors of these two spectra with more data available now and in the future. Also, extending \( E_0 \) to 1.9 GeV or lower, would allow to establish more clearly the stability of the results in this region.

### 7 Conclusions

The \( B \) factories are doing an excellent job of overconstraining the Unitarity Triangle. The extraction of \( |V_{cb}| \) from fits to moments in inclusive \( B \to X_c \ell \bar{\nu} \) decays has reached a better than 2\% precision. \( |V_{cb}| \) from exclusive decays is now known with a 5\% uncertainty.

On the extraction on \( |V_{us}| \), there are encouraging results from Belle, that has calculated branching fractions from exclusive \( B \to X_u \ell \bar{\nu} \). There are also recent theoretical developments that can allow the extraction of \( |V_{us}| \) with a reduced model dependence, exploiting the increasing dataset collected at the \( B \) factories.

Finally, the measurement of relative branching fractions for \( B \to D, D^*, D^{**} \ell \bar{\nu} \) decays and of the form factors in \( B^0 \to D^{*+} \ell^{-} \nu_{\ell} \) decays has improved our modeling of the hadronic system \( X_c \), crucial in the extraction of \( |V_{cb}| \) and in isolating \( B \to X_u \ell \bar{\nu} \) decays for determining \( |V_{ub}| \).
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Figure 5: Comparison of $|V_{ub}|$ values extracted from the three different calculations as a function of the lepton energy cut-off, $E_0$.

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Radiative $b \rightarrow d$ Penguins

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Abstract

This article gives an overview of the recent searches and measurements of $b \rightarrow d$ penguin transitions with the BaBar experiment. The branching fraction of these decays in the Standard Model (SM) is expected to be a factor of 10 or more lower than the corresponding $b \rightarrow s$ penguin transitions, but a deviation from the SM prediction would be an equally striking sign of new physics. The exclusive decay $B \rightarrow \pi \ell \ell$ is searched by BaBar with no excess over the background found. The BaBar measurement of $B \rightarrow (\rho, \omega)\gamma$ provides the first evidence of $B^+ \rightarrow \rho^+ \gamma$,$^c$ is in good agreement with the previous Belle results and provides a measurement of $|V_{td}/V_{ts}|$ independent of the one from $B_s$ mixing. No deviation from the SM is found.

1 Introduction

Since the first $\approx 10$ radiative penguin decays of B mesons have been reported by CLEO in 1993$^1$, the measurements of radiative penguins at the B factory experiments BaBar and Belle have expanded into a rich field of physics. As many as 15 individual exclusive $b \rightarrow s$ modes have been observed and some of the corresponding branching fractions have been measured with a precision of about 10%. This makes it possible to test SM predictions of these rare decays, e.g. QCD calculations of form factors. Additionally, the inclusive measurement of $b \rightarrow s\gamma$ allows strong constraints on new physics models$^2$ and precision measurements of heavy quark effective theory predictions.

With the success of the $b \rightarrow s$ penguin measurements in hand, the next step is to explore $b \rightarrow d$ penguins, which due to Cabibbo suppression are not only about one order of magnitude more rare, but also face a much more severe background due to the abundance of pions in the background. This paper gives an overview of the recent results from BaBar for the exclusive modes $B \rightarrow \pi \ell \ell$ and $B \rightarrow (\rho, \omega)\gamma$. Section 2 will briefly discuss the interest in these measurements from the point of view of the search for new physics, Sections 3 and 4 will summarize the measurements, an discuss the measurement of $|V_{td}/V_{ts}|$ from $B(B \rightarrow (\rho, \omega)\gamma)$.

2 New Physics in Radiative Penguins

The beauty of the search for new physics in penguin decays lies in the fact that the dominating diagram in the SM is a loop or box diagram with heavy particles in the loop. This is exem-
Figure 1: New Physics in Radiative Penguins. The left graph shows the SM contribution, the middle graph a possible SUSY contribution at the same loop level and with similar order of magnitudes of the couplings and masses. The right graph shows a SUSY or 2HDM contribution.

Figure 2: The dominating $b \to d \ell \ell$ penguin graphs in the SM. Similar new physics contributions as in Fig. 1 apply.

plified in Figure 1. Possible new physics, Supersymmetry (SUSY) or the two-Higgs Doublet Model (2HDM) in this case, contribute with particles in the same mass range of 100 to several hundred GeV and often also with similar couplings. This allows strong constraints on new physics models. However, while the measurement of exclusive decays is often experimentally clean, the interpretation of the experimental results for very rare exclusive decays often suffers from theoretical uncertainties in the prediction of form factors, radiative corrections, and other suppressed graphs. Still, strong constraints and precise SM measurements can be gained from ratios of branching fractions.

3 $b \to d \ell^+ \ell^-$ Transitions

The smallest $B$ branching fractions measured up to now are the $b \to s$ modes $B(B \to K \ell \ell) = (3.4 \pm 0.7 \pm 0.2) \times 10^{-7}$ and $B(B \to K^{*} \ell \ell) = (7.8 \pm 1.9 \pm 1.1) \times 10^{-7}$. In comparison to that, $B(B \to \pi \ell \ell)$ is expected to be suppressed by $|V_{td}/V_{ts}|^2$, with the most recent prediction of $B(B \to \pi \ell \ell) = 3.3 \times 10^{-8}$. Any deviation from this order of magnitude would be a striking sign of new physics. The two dominating diagrams for this transition are shown in Figure 2, showing one penguin and one $W$ box diagram.

The BaBar analysis on a dataset of 209 fb$^{-1}$ faces another experimental challenge with respect to $B \to K \ell \ell$ in addition to the reduced branching fraction, namely the much higher rate of $\pi$ in the background than $K$. The analysis strategy is to first select clean $\pi$, $e$ and $\mu$ candidates and then veto resonances decaying into $\ell \ell$. The $u\bar{u}, d\bar{d}, s\bar{s}$ events are strongly reduced by requiring two $p > 1$ GeV leptons in the Event. The charmonium veto against $B \to J/\psi(\psi') \pi(K^*)$ events is shown in Figure 3 (a). The tilt in the mass veto stems from the fact that events with bremsstrahlung of one of the leptons are off both in mass and in reconstructed energy, visible in $\Delta E = E_B - 1/2 E_{Beam}$. Event shape variables against continuum events are grouped in a Fisher discriminant, and event shape variables against $B\bar{B}$ background events are grouped into a likelihood. After these cuts, only combinatoric background from $c\bar{c}$ and $BB$ events are left.

Backgrounds are controlled with various control samples ($e\mu, B \to J/\psi(\psi') \pi(K^*)$). Hadronic mistags are measured in a separate study by inverting the hadron vetoes and then imposing the measured mistag rates as event weights. The resulting small peaking background fraction is shown in Figure 3 (b) with $\pi^+h^+h^-$ on the left and $\pi^0h^+h^-$ on the right side. Finally, the background in the $m_{ES}$ and $\Delta E$ sidebands is extrapolated into the signal region (see Figure 4) to assess the background level independent of the MC simulation.
Figure 3: In (a), the \( m_{\Delta E}, \Delta E \) distribution with the charmonium vetos in the \( B \rightarrow \pi \ell \ell \) analysis is shown. (b) shows the peaking background determination in \( B^0 \rightarrow \pi^0 \ell \ell \) (left) and \( B^+ \rightarrow \pi^+ \ell \ell \) (right).

Figure 4: Results of the \( B \rightarrow \pi \ell \ell \) analysis in the different modes in the \( m_{E_{	ext{CM}}}, \Delta E \) plane. The signal boxes are shown in red.

No excess over the background is observed, hence limits are set using a frequentist cut-and-count method in the signal boxes. The resulting limit of \( B(B \rightarrow \pi \ell \ell) = 3.3 \times 10^{-8} \) assuming isospin symmetry is within a factor of 3 of the SM prediction. A detailed summary of the results can be found in Table 1.

4 \( b \rightarrow d \gamma \) Transitions

\( b \rightarrow d \gamma \) penguins have been first observed by Belle with 350 fb\(^{-1}\) with evidence for the \( B^0 \rightarrow \rho^0 \gamma \) mode\(^6\). In addition to the possibility of finding new physics if the measured branching fractions exceed \( \approx 0.5 \times 10^{-6} \) for the neutral mode or \( \approx 1 \times 10^{-6} \) for the charged mode significantly, it offers the important possibility to measure \( |V_{td}/V_{ts}| \) via

\[
\frac{\Gamma(B \rightarrow \rho \gamma)}{\Gamma(B \rightarrow K^* \gamma)} = \left| \frac{V_{td}}{V_{ts}} \right|^2 \left( \frac{m_B - m_\rho}{m_B - m_{K^*}} \right)^2 \left( \frac{T^{\rho}(0)}{T^{K^*}(0)} \right)^2 (1 + \Delta R),
\]

with \( \Delta R = 0.1 \pm 0.1 \)\(^7\) containing radiative corrections and sub-dominant helicity suppressed WW-diagram (hence depending on \( V_{ts} \) and the CKM fit itself and thus not uncorrelated from \( |V_{td}/V_{ts}| \)), and a form-factor ratio of \( T^{K^*}(0)/T^{\rho}(0) = 1.17 \pm 0.09 \)\(^8\). The graphs involved here are expected to have completely independent possible new physics contributions than \( \Delta m_d/\Delta m_s \), hence the comparison of the two results is very important.
Table 1: Results of the $B \rightarrow \pi \ell \bar{\ell}$ selection and limits on the branching fraction inferred from the absence of an excess over the SM background.

<table>
<thead>
<tr>
<th>Mode</th>
<th>exp.</th>
<th>backg.</th>
<th>90% CL ($10^{-7}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^{\pm} \rightarrow \pi^{\pm} \ell \bar{\ell}$</td>
<td>2.86 ± 0.38</td>
<td>2.17</td>
<td></td>
</tr>
<tr>
<td>$B^0 \rightarrow \pi^0 \ell \bar{\ell}$</td>
<td>0.71 ± 0.30</td>
<td>1.15</td>
<td></td>
</tr>
<tr>
<td>isospin combination</td>
<td></td>
<td></td>
<td>0.91</td>
</tr>
<tr>
<td>$B \rightarrow \pi e \mu$</td>
<td>2.77 ± 0.70</td>
<td>0.92</td>
<td></td>
</tr>
</tbody>
</table>

Figure 5: Extraction of the signal in the $B \rightarrow \rho/\omega \gamma$ analysis in the $B^+ \rightarrow \rho^+ \gamma$ channel in a four-dimensional simultaneous extended maximum likelihood fit. The signal is shown dash-dotted (red), the continuum is dashed (black), the total $B$ background is shown dotted (blue).

Here, the experimental challenge lies not only in the particle ID requirements to suppress $K$ background, but also in the $\pi$ combinatorics coming from the wide resonance states ($\Gamma(\rho) = 150$ MeV) and the high photon background from continuum events with $\pi^0/\eta \rightarrow \gamma \gamma$. The BaBar analysis $^9$ on $316 \text{fb}^{-1}$ of data tackles this challenge by a likelihood veto against $\pi^0/\eta$ based on the invariant mass of the photon pair and the energy of the lower energetic photon, improving the veto significantly over a simple cut on $m_{\gamma\gamma}$. Additionally, a Neural Net (NN) based continuum suppression is applied.

To extract the result, a simultaneous maximum-likelihood fit is performed to $m_{ES}, \Delta E$, the NN output and $\cos \theta_{had}$, where $\theta_{had}$ is the helicity angle of the vector meson $\rho$ or $\omega$, which are transversely polarized in signal events. For $\omega$, the Dalitz angle of the $\omega$ decay is used additionally. An example fit is shown for the $B^+ \rightarrow \rho^+ \gamma$ mode in Figure 5. Many control sample checks are performed. With off-peak data used to control the continuum simulation, and $B \rightarrow K^* \gamma$ and $B \rightarrow D \pi$ used to control resolutions and efficiencies.

The result for the individual modes and different combinations is shown in Table 2. While the $B^0 \rightarrow \omega \gamma$ signal is not yet significant on its own, this result represents the first evidence for $B^+ \rightarrow \rho^+ \gamma$. Figure 6 shows the good agreement of the results with the SM predictions and the agreement with the earlier Belle results. The result for the isospin combination $B \rightarrow \rho/\omega \gamma$ is to be taken with care, however, since the $\omega$ does not belong to the isospin triplet.
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Figure 6: Graphical representation of the results of the Belle and Babar measurements of $b \to d \gamma$ penguins and the comparison with the SM prediction.

Table 2: Results of the BaBar $B \to (\rho/\omega)\gamma$ analysis in terms of extracted branching fraction and significance. The isospin combination is shown in the last line.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Result</th>
<th>Significance</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B(B^+ \to \rho^+\gamma)$</td>
<td>$(1.1^{+0.35}_{-0.23} \pm 0.09) \times 10^{-6}$</td>
<td>3.8 $\sigma$</td>
</tr>
<tr>
<td>$B(B^0 \to \rho^0\gamma)$</td>
<td>$(0.79^{+0.29}_{-0.20} \pm 0.06) \times 10^{-6}$</td>
<td>4.9 $\sigma$</td>
</tr>
<tr>
<td>$B(B^0 \to \omega\gamma)$</td>
<td>$(0.40^{+0.24}_{-0.30} \pm 0.05) \times 10^{-6}$</td>
<td>2.2 $\sigma$</td>
</tr>
</tbody>
</table>

combination of all modes

$B(B \to \rho/\omega\gamma)$ | $(1.25^{+0.25}_{-0.24} \pm 0.09) \times 10^{-6}$ | 6.4 $\sigma$

The results of $B(B \to \rho/\omega\gamma)$ can be used to extract $|V_{td}/V_{ts}|$ using Eq. 1. Using the world average of $B(B \to \rho/\omega\gamma) = (1.25^{+0.25}_{-0.24} \pm 0.09) \times 10^{-6}$, a value of $|V_{td}/V_{ts}|_{\rho/\omega\gamma} = 0.202^{+0.101}_{-0.016} \pm 0.015$ can be extracted, in very good agreement with the result obtained from $\Delta m_d/\Delta m_s$ of $|V_{td}/V_{ts}|_{\Delta m_d/\Delta m_s} = 0.2060 \pm 0.0007^{+0.0081}_{-0.0060}$ from the Tevatron 11. While the precision from $B \to \rho/\omega\gamma$ is not sufficient to compete with $\Delta m_d/\Delta m_s$, it provides an important independent check of possible new physics, due to the different nature of the possible new physics contributions in the two processes. A comparison of the two results with the CKM fit and the imposed constrained in the $\rho, \eta$ plane of the CKM parameterization can be found in Fig. 7.

5 Summary and Outlook

With the extraordinary luminosities of the B-factories Belle and BaBar, the field of strange radiative penguin decays of the B meson has evolved into one of the most important areas of precision measurements in the search of new physics. While it is still unclear whether $B \to \pi\ell\ell$ transitions will be seen in the lifetime of the present B-factories, the rare decay $b \to d\gamma$ has now been seen by Belle and BaBar, and these measurements will possibly evolve towards precision measurements in the same way as the $b \to s\gamma$ decays before. With the anticipated $L_{int} \approx 1$ ab$^{-1}$ of luminosity per B-factory at the end of 2008, a $\approx 10\%$ measurement of the CP asymmetry in $B \to \rho/\omega\gamma$ should be feasible. Another interesting measurement would be the isospin asymmetry $A_I = 2\Gamma(B^0 \to \rho^0\gamma) / \Gamma(B^\pm \to \rho^\pm\gamma) - 1$, which presents a completely new way of obtaining a measurement of the CKM-angle $\gamma$, as outlined in 10 and shown in Fig. 8.

Acknowledgments

The author would like to thank the PEP-II accelerator and its crew for the excellent luminosity delivered to BaBar, and the radiative penguin analysis group in BaBar for lots of support and
very lively discussions.

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Beyond SM, Extradim

C Physics
BASICS OF $D^0$–$D^0$ MIXING

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Complementing the presentations, at this conference, of the first experimental evidence for $D$ mixing found at BaBar and Belle, I discuss the theoretical status of $D$ mixing.

Quasi-Imromptu Talk given at XLIIInd Rencontres de Moriond, Electroweak Interactions and Unified Theories, La Thuile, Italy, March 2007
The highlight of this year’s Moriond conference on electroweak interactions and unified theories arguably was the announcement by BaBar and Belle of experimental evidence for $D^0$–$ar{D}^0$ mixing\(^1,2\), accompanied by an experimental paper\(^3\) and, less than one week after the event, a theoretical analysis\(^4\) – very likely to be followed by many others. As the experimental result came as a surprise to everyone, including the conference organisers, no theoretical talk on the topic had been organised. I was asked to fill the gap and give a quasi-impromptu talk on the theory basics of $D$ mixing, whose written form is presented in these pages. Excellent reviews on the topic can be found in Refs.\(^5,6,7\), and an enlightening reminder of the importance of charm physics in Ref.\(^7\).

In complete analogy to $B$ mixing, $D$ mixing in the SM is due to box diagrams with internal quarks and $W$ bosons. In contrast to $B$, though, the internal quarks are down-type. Also in contrast to $B$ mixing, the GIM mechanism is much more effective, as the heaviest down-type quark, the $b$, comes with a relative enhancement factor $(m_b^2 - m_{s,d}^2)/(m_s^2 - m_d^2)$, but also a large CKM-suppression factor $|V_{ub}V_{c{\bar{b}}}|^2/|V_{us}V_{c{\bar{s}}}|^2 \sim \lambda^5$, which renders its contribution to $D$ mixing $\sim 1\%$ and hence negligible. As a consequence, $D$ mixing is small in the SM, which makes it very sensitive to the potential intervention of new physics (NP), but on the other hand, it is also more difficult to accurately calculate the SM "background", as the loop-diagrams are dominated by $s$ and $d$ quarks and hence sensitive to the intervention of resonances and non-perturbative QCD, see Fig. 1. The quasi-decoupling of the 3rd quark generation also implies that CP violation in $D$ mixing is extremely small in the SM, and hence any observation of CP violation will be a clear-cut signal of new physics, independently of hadronic uncertainties.

The theoretical parameters describing $D$ mixing can be defined in complete analogy to those for $B$ mixing: the time evolution of the $D^0$ system is described by the Schrödinger equation

$$
\frac{\partial}{\partial t} \begin{pmatrix} D^0(t) \\ \bar{D}^0(t) \end{pmatrix} = -i \left( M - i \frac{\Gamma}{2} \right) \begin{pmatrix} D^0(t) \\ \bar{D}^0(t) \end{pmatrix}
$$

with Hermitian matrices $M$ and $\Gamma$. The off-diagonal elements of these matrices, $M_{12}$ and $\Gamma_{12}$, describe, respectively, the dispersive and absorptive parts of $D$ mixing. The flavour-eigenstates $D^0 = (c\bar{u})$, $\bar{D}^0 = (u\bar{c})$ differ from the mass-eigenstates $D_{1,2}$; they are related by

$$
|D_{1,2}\rangle = p|D^0\rangle \pm q|\bar{D}^0\rangle
$$

with

$$
|q|^2 = \frac{M_{12}^2 - \frac{1}{2} \Gamma_{12}^2}{M_{12}^2 - \frac{1}{2} \Gamma_{12}^2},
$$

$$
|p|^2 = \frac{M_{12}^2 - \frac{1}{2} \Gamma_{12}^2}{M_{12}^2 - \frac{1}{2} \Gamma_{12}^2}.
$$

Figure 1: Resonance contribution to $D$ mixing. Figure taken from Ref.\(^5\).
The basic observables in $D$ mixing are the mass and lifetime differences of $D_{1,2}$, which are usually normalised to the average lifetime $\Gamma = (\Gamma_1 + \Gamma_2)/2$:

$$\Delta M/\Gamma = \frac{M_2 - M_1}{\Gamma}, \quad \Delta \Gamma/2\Gamma = \frac{\Gamma_2 - \Gamma_1}{2\Gamma}. \quad (4)$$

While previously only bounds on $x$ and $y$ were known, both BaBar and Belle have now obtained evidence for a non-vanishing mixing in the $D$ system. BaBar has obtained this evidence from the measurement of the doubly Cabibbo-suppressed decay $D^0 \rightarrow K^+\pi^-$ (and its CP conjugate), yielding

$$y' = (0.97 \pm 0.44(\text{stat}) \pm 0.31(\text{syst})) \times 10^{-2}, \quad x'^2 = (-0.022 \pm 0.030(\text{stat}) \pm 0.021(\text{syst})) \times 10^{-2},$$

while Belle obtains

$$y_{\text{CP}} = (1.31 \pm 0.32(\text{stat}) \pm 0.25(\text{syst})) \times 10^{-2} \quad (5)$$

from $D^0 \rightarrow K^+K^-, \pi^+\pi^-$ and

$$x = (0.80 \pm 0.29(\text{stat}) \pm 0.17(\text{syst})) \times 10^{-2}, \quad y = (0.33 \pm 0.24(\text{stat}) \pm 0.15(\text{syst})) \times 10^{-2} \quad (7)$$

from a Dalitz-plot analysis of $D^0 \rightarrow K^0_S\pi^+\pi^-$. Here $y_{\text{CP}} \rightarrow y$ in the limit of no CP violation in $D$ mixing, while the primed quantities $x', y'$ are related to $x, y$ by a rotation by a strong phase $\delta$, see below.

CP violation in $D^0 \rightarrow f$ decays, which is predicted to be extremely small in the SM, can be characterised by non-vanishing values of

$$A_M = \left| \frac{q}{p} \right| - 1, \quad \phi = \arg(M_{12}/\Gamma_{12}),$$

where $A_M$ measures CP violation in the mixing amplitude, while $\phi$ plays a rôle in the interference between the decays $D^0 \rightarrow f$ and $\bar{D}^0 \rightarrow f$.

Various $D^0$ decay channels are sensitive to $D$ mixing. The evidence found by BaBar relies on $D^0 \rightarrow K^+\pi^-, \bar{D}^0 \rightarrow K^-\pi^+$, which are “wrong sign” decays (the dominant transition is $c \rightarrow s$ and produces a $K^-$ in $D^0$, and a $K^+$ in $\bar{D}^0$ decays) and receive contributions from two amplitudes: a doubly Cabibbo-suppressed amplitude $\bar{D}^0 \rightarrow K^-\pi^+$, i.e. $\bar{c} \rightarrow d\bar{u}s$, and a two-step process via the oscillation $\bar{D}^0 \rightarrow D^0$, followed by the Cabibbo-favoured process $D^0 \rightarrow K^-\pi^+$, see Fig. 2. The relevant point here is that the amplitude with no oscillation is heavily suppressed which makes it competitive with the oscillated amplitude. The wrong-sign time-dependent
decay rate $D^0(t) \to K^+\pi^-$ is usually normalised to the Cabbibo-favoured rate $D^0 \to K^-\pi^+$. Expanding the ratio of suppressed vs. favoured amplitudes to second order in $x, y$, one finds

$$\frac{\Gamma(D^0(t) \to K^+\pi^-)}{\Gamma(D^0 \to K^-\pi^-)} = \Gamma e^{-\Gamma t} \left[ R_D + \frac{q}{p} \sqrt{R_D(y'\cos\phi - x'\sin\phi)}(\Gamma t) + \frac{q^2 y^2 + y'^2}{4} (\Gamma t)^2 \right]$$

$$\frac{\Gamma(D^0(t) \to K^-\pi^+)}{\Gamma(D^0 \to K^-\pi^-)} = \Gamma e^{-\Gamma t} \left[ R_D + \frac{p}{q} \sqrt{R_D(y'\cos\phi + x'\sin\phi)}(\Gamma t) + \frac{p^2 x^2 + y^2}{4} (\Gamma t)^2 \right]$$

(9)

where the overall factor $\Gamma$ ensures the correct normalisation upon integration over $t$. Here $R_D^{1/2}$ is the modulus of the ratio of the doubly Cabbibo-suppressed amplitude vs. the favoured one and $x', y'$ contain the effect of the relative strong phase $\delta$ between the two amplitudes:

$$A(D^0 \to K^+\pi^-) = -R_D^{1/2} e^{-i\delta},$$

$$x' = x \cos \delta + y \sin \delta, \quad y' = y \cos \delta - x \sin \delta;$$

(10)

(11)

$\delta$ vanishes in the SU(3) limit. The minus sign in (10) originates from the sign of $V_{us}$ relative to $V_{ud}$. Note that the 2nd term in brackets in (9) comes from the interference of the two decay amplitudes with and without mixing. BaBar has obtained $R_D, y'$ and $x^2$ from the fit of their experimental results to the above formulas in the case of (a) CP conservation, i.e. $|q/p| \to 1, \phi \to 0$, and (b) CP violation, i.e. different coefficients $R_D^{1/2}, y'^2$, etc. for $D^0$ and $\bar{D}^0$ decays. The difference of the latter proved to be compatible with 0, so there is no evidence for CP violation.

Let us now turn to the theoretical predictions for $x$ and $y$ in the SM. In terms of hadronic matrix elements, $M_{12} = M_{21}^{\text{ph}}$ and $\Gamma_{12} = \Gamma_{21}^{\text{ph}}$ can be expressed as

$$M_{12} = \langle D^0 | \mathcal{H}_{\text{eff}}^{C-2} | D^0 \rangle + P \sum_n \frac{\langle D^0 | \mathcal{H}_{\text{eff}}^{C-1} | n \rangle \langle n | \mathcal{H}_{\text{eff}}^{C-1} | D^0 \rangle}{m_n^2 - E_n^2},$$

$$\Gamma_{12} = P \sum_n \beta_n^{\text{ph}} \langle D^0 | \mathcal{H}_{\text{eff}}^{C-1} | n \rangle \langle n | \mathcal{H}_{\text{eff}}^{C-1} | D^0 \rangle.$$

(12)

While the expression for $\Gamma_{12}$ is very similar to that in the $B$ system, that for $M_{12}$ differs by the contribution of the second term which is heavily suppressed in $B$ mixing. The sum runs over all decay channels of $D^0$; the contribution to $M_{12}$ includes that of off-shell intermediate states, while only on-shell states contribute to $\Gamma_{12}$; $\beta_n^{\text{ph}}$ is the corresponding phase-space factor. $\mathcal{H}_{\text{eff}}^{C-2}$ is the local Hamiltonian obtained from the box diagrams, and includes potential contributions from NP, while all terms in $\mathcal{H}_{\text{eff}}^{C-1}$, the Hamiltonian describing non-leptonic decays of the $c$ quark, are dominated by SM contributions (see, however, Ref. 8 for a discussion of NP effects in decay amplitudes). Neglecting long-distance non-perturbative QCD effects, and only including the box diagrams, one finds $x_{\text{box}} = O(10^{-5}), y_{\text{box}} = O(10^{-7})$, which is far below the experimental results – which indicates that these long-distance effects are extremely important.

There is an extensive literature on estimating $x$ and $y$ within and beyond the SM, see Ref. 10 for a collection of results. The central problem of all these calculations is that the $D$ is too heavy to be treated as light and too light to be treated as heavy. As a consequence, the two approaches that have been so successful in treating heavy ($B$) and light ($K$) meson mixing both are not really applicable to $D$ mixing: the "inclusive" approach is based on operator product expansion and relies on quark-hadron duality. If $\Lambda/m_c$, where $\Lambda$ is a hadronic scale, is considered a small parameter, $x$ and $y$ can be expanded in terms of matrix elements of local operators.
series can be truncated after a few terms. Such calculations typically yield $x, y \lesssim 10^{-3}$, and the result of both BaBar and Belle, $y \sim 10^{-2}$, is certainly not a generic prediction of such an analysis. In the "exclusive" approach\textsuperscript{9,12}, on the other hand, one sums over intermediate hadronic states, which may be modeled or fit to experimental data. One crucial observation\textsuperscript{9} is that $x$ and $y$ are only generated at second order in SU(3) breaking, which suggests an analysis based on the summation over exclusive states arranged in SU(3) multiplets. As argued in Ref.\textsuperscript{9}, the main source of SU(3)-breaking within these multiplets is due to phase-space, or rather, the lack thereof: if the heaviest members of a multiplet are too heavy to be kinematically accessible in the decay, they have to be excluded from the sum over all members of the multiplet (e.g. $D \to 4\pi$ is kinematically allowed, but $D \to 4K$ is not) and as a consequence, the cancellation of the sum over all terms, which yields $0$ in the SU(3)-limit, is badly broken. The conclusion is that in this way values of $y \sim 10^{-2}$ can be reached – which agrees very well with the experimental result and suggests that these threshold effects may indeed explain the experimental result. The inclusive approach, on the other hand, relies on the duality of hadronic and partonic effects, smeared over sufficiently large energy intervals, and is manifestly insensitive to threshold phenomena – and hence likely to be inapplicable to $D$ decays. In the exclusive approach, $x$ can be related to $y$ via a dispersion relation; the authors of Ref.\textsuperscript{9} find that for $y \sim 1\%$ one expects $|x|$ between 0.1% and 1%, and $x$ and $y$ to be of opposite sign; one should be aware, however, that this calculation is more model-dependent than that of $y$.

In conclusion, we find that the experimental results on $D$ mixing reported by BaBar and Belle at the 2007 Rencontres de Moriond on electroweak interactions and unified theories present a major step forward in experimental achievement and analysis. The measured value of $y > x$ is at the high end of theoretical predictions and indicates large long-distance contributions, which also impact on $x$, i.e. the short-distance/NP sensitive mass difference. As long as there is no major breakthrough in theoretical predictions for $D$ mixing, which are held back by the fact that the $D$ meson is at the same time too heavy and too light for our current theoretical tools to get a proper grip on the problem, the long-distance SM contributions to $x$ will completely obscure any NP contributions and their detection. The observation of CP violation still presents a theoretically clean way for NP to manifest itself and it is to be hoped that in the near future, i.e. at the $B$ factories or the LHC, at least one of the plentiful opportunities for NP to show up in CP violation\textsuperscript{13} will be realised.

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1. M. Staric, talk given at this conference.
2. K. Flood, talk given at this conference.
3. B. Aubert et al. [BABAR Collaboration], arXiv:hep-ex/0703020.
SEARCH FOR $D^0 - \bar{D}^0$ MIXING AT BELLE AND BABAR

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We discuss recent measurements of the Belle and BaBar collaborations performed to search for mixing in the neutral $D$-meson system. The data were obtained in $e^+e^-$ collisions at or near the $\Upsilon(4S)$ resonance. The analyses cover semi-leptonic decays, wrong-sign hadronic decays, self-conjugate decays and decays to $CP$ eigenstates. In the latter case the Belle collaboration has found the first experimental evidence for the $D$-mixing.

1 Introduction

The phenomenon of mixing between a particle and its anti-particle has been observed in several systems of neutral mesons: neutral kaons, $B^0_d$, and most recently $B^0_s$ mesons. This process is also possible in the $D$-meson system, but has not previously been observed.

The time evolution of a $D^0$ or $\bar{D}^0$ is governed by the mixing parameters $x = (M_1 - M_2)/\Gamma$ and $y = (\Gamma_1 - \Gamma_2)/2\Gamma$, where $M_{1,2}$ and $\Gamma_{1,2}$ are the masses and widths, respectively, of the mass eigenstates, and $\Gamma = (\Gamma_1 + \Gamma_2)/2$. For no mixing, $x = y = 0$. Within the Standard Model (SM) the rate of $D$-mixing is expected to be small due to the near degeneracy of the $s$ and $d$ quark masses relative to the $W$ mass, and the small value of the $b$ quark couplings. Predictions for $x$ and $y$ are dominated by non-perturbative processes that are difficult to calculate. The largest predictions are $|x|, |y| \sim O(10^{-3})$. Loop diagrams including new, as-yet-unobserved particles could significantly affect the experimental values. $CP$-violating effects in $D$-mixing would be a clear signal of new physics, as $CP$ violation ($CPV$) is expected to be very small in the SM.

Both semi-leptonic and hadronic $D^0$ decays have been used to constrain $x$ and $y$. In order to tag the flavour at the production the $D^0$ meson is usually reconstructed in the decay $D^{*+} \rightarrow D^0\pi^+_s$, where the charge of a characteristic slow pion $\pi_s$ tags the initial $D^0$ flavour. Usually also the $D^0$ proper decay time is measured, since the decay time distribution of mixed

$^a$Charge conjugate modes are implied unless explicitly stated otherwise.
events depends on the mixing parameters $x$ and $y$ and differs from that of un-mixed events. The proper decay time of an event is determined from the distance between the production and the decay vertex. The decay vertex is obtained from $D^0$ daughter tracks, refitted to originate from a common point. The production vertex is found by constraining the $D^0$ momentum vector to originate from the $e^+e^-$ interaction region. The proper decay time resolution is typically equal to a half of the $D^0$ lifetime.

The measurements are performed at $\Upsilon(4S)$ resonance. To reject $D^{*+}$ mesons coming from the decays of $B$ mesons a cut on the $D^{*+}$ or $D^0$ momentum in the $e^+e^-$ center-of-mass frame (CM) is applied at about 2.5 GeV/c. This cut also suppresses the combinatorial background.

2 Semileptonic decays

In semileptonic decays the wrong-sign (WS) process $D^{*+} \to K^+ \ell^- \nu \pi^+_\ell$ can proceed only through $D^0$ mixing followed by the right-sign (RS) decay: $D^0 \to \overline{D}^0 \to K^+ \ell^- \nu$. The mixing rate $R_M = \frac{1}{2}(x^2 + y^2)$ can be measured in the time-integrated way. By counting WS and RS events ($N_{WS}, N_{RS}$) and by applying small efficiency corrections ($\epsilon_{WS}, \epsilon_{RS}$) one can determine the mixing rate with

$$R_M = \frac{N_{WS}}{N_{RS}} \frac{\epsilon_{RS}}{\epsilon_{WS}} \quad (1)$$

The Belle collaboration published their measurement in semileptonic decays with an electron in the final state using 253 fb$^{-1}$ of data. They reconstructed 229 $\times 10^3$ RS signal events with a purity of 73%. In the WS spectra they found no evidence for a signal. They reported an upper limit for the mixing rate of $R_M < 1.2 \times 10^{-3}$ @ 95% C.L.

The new measurement by BaBar collaboration performed on 344 fb$^{-1}$ of data uses a double flavor tagging method for background suppression. With this method the initial flavor is tagged twice: once using the slow pion and once using the flavor of a fully reconstructed $D$ meson in the hemisphere opposite to the semileptonic candidate. Both tags are required to be consistent. After applying their final selection criteria, they found 4780 RS signal events and 3 WS candidates. The expected WS background was 2.85 events. The upper limit for the mixing rate obtained from these numbers is $R_M < 1.2 \times 10^{-3}$ @ 90% C.L.

3 Wrong-sign decay $D^0 \to K^+\pi^-$

The WS final state $K^+\pi^-$ can be accessed either directly via doubly Cabibbo-suppressed (DCS) decay or through mixing followed by the Cabibbo-favored (CF) decay $D^0 \to \overline{D}^0 \to K^+\pi^-$. The two decays can be distinguished by the decay-time distribution. For $|x|, |y| \ll 1$, and assuming negligible $CP$ violation, the decay-time distribution for the WS process can be expressed as

$$\frac{dN}{dt} \propto |R_D + y' \sqrt{R_D} (\Gamma_t) + \frac{x'^2 + y'^2}{4} (\Gamma_t)^2 | e^{-\Gamma t}$$

where $R_D$ is the ratio of DCS to CF decay rates, $x' = x \cos \delta + y \sin \delta$, $y' = y \cos \delta - x \sin \delta$, and $\delta$ is the strong phase difference between DCS and CF amplitudes. The first (last) term is due to the DCS (CF) amplitude, and the middle term is due to interference between the two processes.

The Belle collaboration reported the results of a search for $D$-mixing in this decay mode based on 400 fb$^{-1}$. They found $1.07 \times 10^6$ RS and 4024 WS signal events with a ratio of WS signal to background of 1.1. The RS events were used to establish the resolution function parameters and the $D^0$ lifetime, which they found to be consistent with the world average. The unbinned maximum likelihood fit of the WS data assuming $CP$ conservation resulted in $R_D = (0.364 \pm 0.017)\%$, while the mixing parameters $x'^2 = (0.15^{+0.21}_{-0.32}) \times 10^{-3}$ and $y' = (0.6^{+4.0}_{-3.3}) \times 10^{-3}$ are consistent with zero. They fitted also with the assumption of $CP$ violation; again the mixing
parameters and the CP asymmetries in mixing and in decay are consistent with zero. The upper limit for the mixing rate is $R_M < 0.40 \times 10^{-3}$ @ 95% C.L.

Fig. 2 shows the 95% C.L. contours. In the case of CP conservation, the no-mixing point $x^2 = y^2 = 0$ lies just outside the 95% C.L. contour; this point corresponds to a C.L. of 3.9%. The BaBar collaboration reported just after this presentation a positive $D$-mixing result in this decay mode and on a similar-size data set; the no-mixing point in their measurement corresponds to a C.L. of $1 \times 10^{-4}$.

4 Wrong-sign decays $D^0 \rightarrow K^+\pi^-\pi^0$ and $D^0 \rightarrow K^+\pi^-\pi^+\pi^-$

As in the $K^+\pi^-$ final state, the WS decays $D^0 \rightarrow K^+\pi^-\pi^0$ and $D^0 \rightarrow K^+\pi^-\pi^+\pi^-$ can proceed either through DCS decay or through mixing followed by CF decay. For a nonleptonic multibody WS decay, the decay-time distribution is approximated by

$$\frac{dN}{dt} \propto |\bar{R}_D + \alpha \gamma \sqrt{R_D(t)} \left( \frac{\bar{x}^2 + \bar{y}^2}{4} \right) | \gamma t | e^{-\gamma t}, \quad 0 \leq \alpha \leq 1,$$

similar to Eq. 2. The tilde indicates that the corresponding quantities from Eq. 2 have been integrated over a certain region of phase-space. The parameter $\alpha$ is a suppression factor that accounts for strong-phase variation over the phase-space. The time-integrated mixing rate is $R_M = \frac{1}{2} (x^2 + y^2) = \frac{1}{2} (\bar{x}^2 + \bar{y}^2)$ and is independent of the position in the Dalitz plot.

The BaBar collaboration analyzed 230 fb$^{-1}$ of data to search for $D$-mixing in both decay modes. In the $K^+\pi^0$ analysis they reduced the ratio of DCS to CF decays by excluding events with two-body invariant masses in the range of the $K^{*+}$ and $K^{*0}$ resonances. Their final selection included 503 $\times$ 10$^3$ RS and 763 WS signal events. Assuming CP conservation they found that the data were consistent with the no-mixing hypothesis at the 4.5% confidence level, and they set an upper limit $R_M < 0.54 \times 10^{-3}$ @ 95% C.L.

In the $K^+\pi^+$ analysis no phase-space selection was made. The signal yields were $700 \times 10^3$ events for the RS and 2202 events for the WS final states. Assuming CP conservation, the data were consistent with no mixing at the 4.3% confidence level. The upper limit was set to $R_M < 0.48 \times 10^{-3}$ @ 95% C.L. By combining the results of both analyses they determined the upper limit $R_M < 0.42 \times 10^{-3}$ @ 95% C.L. for the time-integrated mixing rate. The no-mixing hypothesis was found to be at the 2.1% C.L.
5 Time-dependent Dalitz analysis of $D^0 \rightarrow K^0_\pi^+\pi^-$

A new measurement in the self-conjugate decays $D^0 \rightarrow K^0_\pi^+\pi^-$ based on 540 fb$^{-1}$ was performed by the Belle collaboration using a time-dependent Dalitz plot analysis. The time dependence of the $K^0_\pi^+\pi^-$ Dalitz plot distribution allows one to measure $x$ and $y$ directly. This method was developed by CLEO using 9.0 fb$^{-1}$.

Assuming CP conservation, the decay amplitude at time $t$ of an initially produced $D^0$ or $\bar{D}^0$ can be expressed as

$$
\mathcal{M}(m^2_-, m^2_+, t) = A(m^2_-, m^2_+) \frac{e_1(t) + e_2(t)}{2} + A(m^2_-, m^2_-) \frac{e_1(t) - e_2(t)}{2}
$$

where $A$ is the decay amplitude as a function of the invariant masses squared $m^2_\pm$, defined with the $D^*$ tag

$$
m^2_\pm = \begin{cases} 
    m(K_\pi^+\pi^-)^2 & D^{*+} \rightarrow D^0\pi^+ \\
    m(K_\pi^-\pi^+)^2 & D^{*-} \rightarrow D^0\pi^-
\end{cases}
$$

The time dependence is contained in the terms $e_{1,2}(t) = \exp[-i(m_{1,2} - i\Gamma_{1,2}/2)t]$. Upon squaring $\mathcal{M}$, one obtains decay rates containing terms $\exp(-\Gamma_t \cos(x\Gamma_t))$, $\exp(-\Gamma_t \sin(x\Gamma_t))$ and $\exp[-(1 + y)\Gamma_t]$.

The overall decay amplitude $A$ can be expressed as a sum of quasi-two-body amplitudes $A_\nu$ and a constant non-resonant term (NR):

$$
A(m^2_-, m^2_+) = \sum_\nu A_\nu e^{i\phi_\nu} A_\nu(m^2_-, m^2_+) + a_{NR} e^{i\phi_{NR}}
$$

The functions $A_\nu$ are products of Blatt-Weiskopf form factors and relativistic Breit-Wigner functions.

The $D^0$ candidates were reconstructed in the decay chain $D^{*+} \rightarrow D^0\pi^+$, $D^0 \rightarrow K^0_\pi^+\pi^-$ and the charge of $\pi_\pm$ was used to tag the $D^0$ flavor at production. The $K^0_\pi^+$ was reconstructed in the decay to the $\pi^+\pi^-$ final state; an invariant mass within ±10 MeV of $m_K\pi$ and a common vertex separated from the interaction region were required. The $D^0$ decay point was constructed from charged pion tracks only and the production point was obtained from the intersection of the $D^0$ momentum vector with the $e^+e^-$ interaction region.

The signal and background yields were determined from a two-dimensional fit to the variables $M \equiv m_K\pi\pi$ and $Q \equiv m_K\pi\pi - m_\pi$. The background was classified into two types: random $\pi_\pm$ background and combinatorial background. In the signal region, defined as $3\sigma$ intervals in $M$ and $Q$, they found $534 \times 10^3$ signal events and background fractions of 1% and 4% for the random $\pi_\pm$ and combinatorial backgrounds, respectively.

For the events in the signal region they performed a simultaneous unbinned likelihood fit to the Dalitz plot variables $m^2_\pm$ and $m^2_\pi$, and the decay time $t$. The likelihood function is:

$$
\mathcal{L} = \sum_{i=1}^{N} \sum_j f_j(M_i, Q_i) P_j(m^2_-, m^2_+, t_i)
$$

where $j = \{sig, rnd, cmb\}$ denotes the signal or background components, and the index $i$ runs over all events. The event weights $f_j$ are functions of $M$ and $Q$ and were obtained from the $M - Q$ fit mentioned above.

The probability density function (PDF) $P_{sig}$ equals $|M|^2$ convolved with the detector response. The resolution in the decay time $t$ is parameterized by a sum of three Gaussians with a common mean and widths $\sigma_k = S_k \cdot \sigma_t$, ($k = 1, 2, 3$), where $\sigma_t$ is the decay time uncertainty calculated event-by-event, and the $S_k$ are scale factors, left as free parameters in the fit.
The random $\pi_3$ background contains real $D^0$ and $\bar{D}^0$ decays uncorrelated to the charge of a slow pion. The PDF is taken to be $(1 - f_w)|\mathcal{M}(m_{\pi^+}, m_{\pi^0}, t)|^2 + f_w|\mathcal{M}(m_{\pi^0}, m_{\pi^+}, t)|^2$, convolved with the same resolution function as the signal PDF. The fraction $f_w$ was determined from fitting events in the $Q$ sideband.

For the combinatorial background, $P_{\text{comb}}$ is a product of the Dalitz plot and decay time PDFs. The latter is parameterized as the sum of a delta function and an exponential function convolved with a Gaussian resolution function. The resolution function has a $\sigma_t$ dependent offset. This and other timing parameters, as well as the Dalitz PDF, were obtained from the events in the mass sideband.

The free parameters in the fit were $x$, $y$, $\tau_{D^0}$, the timing resolution parameters of the signal, and the Dalitz plot resonance parameters $\sigma_{(N_R)}$ and $\phi_{(N_R)}$. The resonance model assumed 18 quasi-two-body resonances; masses and widths were taken from world averages.

The Dalitz plot and its projections, along with projections of the resulting fit, are shown in Fig. ???. The decay-time distribution and the fit are shown in Fig. ???. The fitted $D^0$ lifetime $\tau_{D^0} = 409.9\pm1.0$ fs is consistent with the world average. The preliminary results for the mixing parameters are $x = (0.80\pm0.29\pm0.17)\%$ and $y = (0.33\pm0.24\pm0.15)\%$. The largest contributions to the systematic uncertainty of the result are found to arise from model dependence and from the fit to the decay-time distribution. The result for the mixing parameter $x$ represents the most stringent limit on this parameter obtained up to now.

6 Decays to $CP$ eigenstates $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$

Belle collaboration has performed a new measurement in the decays to $CP$ eigenstates $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$; treating the decay time distributions as exponential, they measured the quantity

$$y_{CP} = \frac{\tau(K^-\pi^+)}{\tau(K^+\pi^-)} - 1,$$

where $\tau(K^+K^-)$ and $\tau(K^-\pi^+)$ are the lifetimes of $D^0 \to K^+K^-$ (or $\pi^+\pi^-$) and $D^0 \to K^-\pi^+$ decays. It can be shown that $y_{CP} = y\cos\phi - \frac{1}{2}A_M x \sin\phi$, where $A_M$ parameterizes $CPV$ in
mixing and $\phi$ is a weak phase. If $CP$ is conserved, $A_M = \phi = 0$ and $y_{CP} = y$. To date several measurements of $y_{CP}$ have been reported; the average value is $\sim 2$ standard deviations ($\sigma$) above zero. The new Belle measurement, based on 540 fb$^{-1}$, yields a nonzero value of $y_{CP}$ with $> 3\sigma$ significance. They also searched for $CPV$ by measuring the quantity

$$A_\Gamma = \frac{\tau(D^0 \rightarrow K^- K^+)}{\tau(D^0 \rightarrow K^- K^+)} + \frac{\tau(D^0 \rightarrow K^+ K^-)}{\tau(D^0 \rightarrow K^- K^+)};$$

(9)

this observable equals $A_\Gamma = \frac{1}{2} A_M y \cos \phi - x \sin \phi$.

They reconstructed $D^{\ast+} \rightarrow D^0 \pi^+_s$ decays, and $D^0 \rightarrow K^+ K^-, K^- \pi^+$, and $\pi^+ \pi^-$. $D^0$ daughter tracks were refitted to a common vertex, and the $D^0$ production vertex was found by constraining its momentum vector and the $\pi_s$ track to originate from the $e^+e^-$ interaction region. A $D^{\ast+}$ momentum greater than 2.5 GeV/c (in the CM) was required to reject $D$-mesons produced in $B$-meson decays and to suppress combinatorial background. The proper decay time of the $D^0$ candidate was then calculated from the projection of the vector joining the two vertices, $\vec{L}$, onto the $D^0$ momentum vector, $t = m_{D^0} \vec{L} \cdot \vec{p}/p^2$, where $m_{D^0}$ is the nominal $D^0$ mass. The decay time uncertainty $\sigma_t$ was evaluated event-by-event from the covariance matrices of the production and decay vertices.

Candidate $D^0$ mesons were selected using two kinematic observables: the invariant mass of the $D^0$ decay products, $M$, and the energy released in the $D^{\ast+}$ decay, $q = (M_{D^\ast} - M - m_{\pi}) c^2$. $M_{D^\ast}$ is the invariant mass of the $D^0\pi_s$ combination and $m_{\pi}$ is the $\pi^+$ mass.

According to Monte Carlo (MC) simulated distributions of $t$, $M$, and $q$, background events fall into four categories: (1) combinatorial, with zero apparent lifetime; (2) true $D^0$ mesons combined with random slow pions (this has the same apparent lifetime as the signal) (3) $D^0$ decays to three or more particles, and (4) other charm hadron decays. The apparent lifetime of the latter two categories is 10–30% larger than $\tau_{D^0}$.

The sample of events for the lifetime measurements was selected using $|\Delta M|/\sigma_M$, where $\Delta M = M - m_{D^0}$; $|\Delta q| = q - (M_{D^0} - M - m_{\pi}) c^2$; and $\sigma_t$. The invariant mass resolution $\sigma_M$ varies from 5.5–6.8 MeV/$c^2$, depending on the decay channel. Selection criteria were chosen to minimize the expected statistical error on $y_{CP}$, using the MC: they required $|\Delta M|/\sigma_M < 2.3$, $|\Delta q| < 0.80$ MeV, and $\sigma_t < 370$ fs. They found $111 \times 10^3 K^+ K^-$, $1.22 \times 10^6 K^-\pi^+$, and $49 \times 10^3 \pi^+\pi^- \pi^+\pi^-$ signal events, with purities of 98%, 99%, and 92% respectively.

The relative lifetime difference $y_{CP}$ was determined from $D^0 \rightarrow K^+ K^-$, $K^-\pi^+$, and $\pi^+\pi^-$ decay time distributions by performing a simultaneous binned maximum likelihood fit to the three samples. Each distribution was assumed to be a sum of signal and background contributions, with the signal contribution being a convolution of an exponential and a detector resolution function,

$$\frac{dN}{dt} = \frac{N_{sig}}{\tau} \int e^{-\tau/t'} \cdot R(t - t') \, dt' + B(t).$$

(10)
Figure 5: Results of the simultaneous fit to decay time distributions of (a) $D^0 \rightarrow K^+K^-$, (b) $D^0 \rightarrow K^-\pi^+$ and (c) $D^0 \rightarrow \pi^+\pi^-$ decays. The cross-hatched area represents background contributions, the shape of which was fitted using $M$ sideband events. (d) Ratio of decay time distributions between $D^0 \rightarrow K^+K^-, \pi^+\pi^-$ and $D^0 \rightarrow K^-\pi^+$ decays. The solid line is a fit to the data points.

The resolution function $R(t - t')$ was constructed from the normalized distribution of the decay time uncertainties $\sigma_t$ (see Fig. ??). The $\sigma_t$ of a reconstructed event ideally represents an uncertainty with a Gaussian probability density: in this case, bin $i$ in the $\sigma_t$ distribution is taken to correspond to a Gaussian resolution term of width $\sigma_i$, with a weight given by the fraction $f_i$ of events in that bin. However, the distribution of "pulls", i.e. the normalized residuals $(t_{\text{rec}} - t_{\text{gen}})/\sigma_t$ (where $t_{\text{rec}}$ and $t_{\text{gen}}$ are reconstructed and generated MC decay times), is not well-described by a Gaussian. They found that this distribution can be fitted with a sum of three Gaussians of different widths $\sigma_k^{\text{pull}}$ and fractions $w_k$, constrained to the same mean. They therefore choose a parameterization

$$R(t - t') = \sum_{i=1}^{n} f_i \sum_{k=1}^{3} w_k G(t - t'; \sigma_{sk}, t_0),$$

with $\sigma_{sk} = s_k \sigma_k^{\text{pull}} \sigma_t$, where the $s_k$ are three scale factors introduced to account for differences between the simulated and real $\sigma_k^{\text{pull}}$, and $t_0$ allows for a (common) offset of the Gaussian terms from zero.

The background $B(t)$ was parameterized assuming two lifetime components: an exponential and a $\delta$ function, each convolved with corresponding resolution functions as parameterized by Eq. (??). Separate $B(t)$ parameters for each final state were determined by fits to the $t$ distributions of events in $M$ sidebands. The MC was used to select the sideband region that best reproduces the timing distribution of background events in the signal region.

Fits to the $D^0 \rightarrow K^+K^-, K^-\pi^+$ and $\pi^+\pi^-$ data are shown in Fig. ??(a)-(c). The fitted lifetime of $D^0$ mesons in the $K^-\pi^+$ final state, $\tau_{D^0} = (408.7\pm0.6 \text{ (stat.)})$ fs, is in good agreement with the current world average. The result for the apparent lifetime of $D^0$ mesons between decays to $CP$-even eigenstates and the $K^-\pi^+$ final state is

$$\tau_{CP} = (1.31\pm0.32 \text{ (stat.)})\pm0.25 \text{ (syst.)} \%.$$
Combining the errors in quadrature this result is a 3.2 standard deviations from zero and represents the first experimental evidence for the $D$-mixing, regardless of possible CPV. The effect is presented visually in Fig. ??(d), which shows the ratio of decay time distributions for $D^0 \rightarrow K^+K^-, \pi^+\pi^-$ and $D^0 \rightarrow K^-\pi^+$ decays. They also searched for CP violation by separately measuring decay times of $D^0$ and $\bar{D}^0$ mesons in CP-even final states. They find an asymmetry consistent with zero, $A_T = (0.01 \pm 0.30(\text{stat.}) \pm 0.15(\text{syst.}))\%$.

References

8. B. Aubert et al. (BaBar Collaboration), arXiv:hep-ex/0705.0704.
Evidence for $D^0 - \bar{D}^0$ Mixing at Babar

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We present evidence for $D^0 - \bar{D}^0$ mixing in $D^0 \rightarrow K^+\pi^-$ decays from 384 fb\(^{-1}\) of $e^+e^-$ colliding-beam data recorded near $\sqrt{s} = 10.6$ GeV with the BABAR detector at the PEP-II storage rings at SLAC. We find the mixing parameters $x^2 = -0.22 \pm 0.30$ (stat.) $\pm 0.21$ (syst.) $\times 10^{-2}$ and $y^2 = 0.7 \pm 4.4$ (stat.) $\pm 3.1$ (syst.) $\times 10^{-2}$, and a correlation between them of $-0.94$. This result is inconsistent with the no-mixing hypothesis with a significance of 3.9 standard deviations. We measure $R_D$, the ratio of doubly Cabibbo-suppressed to Cabibbo-favored decay rates, to be $0.303 \pm 0.016$ (stat.) $\pm 0.010$ (syst.). We find no evidence for $C\bar{P}$ violation.

1 Introduction

The $D^0$ and $\bar{D}^0$ mesons are flavor eigenstates which are invariant in strong interactions, but are subject to electroweak interactions that permit an initial flavor eigenstate to evolve into a time-dependent mixture of $D^0$ and $\bar{D}^0$. In the Standard Model (SM), such oscillations proceed through both short-distance and non-perturbative long-distance amplitudes. The expected mixing rate mediated by down-type quark box diagrams\(^1\) and chi-penguin\(^2\) diagrams is $\mathcal{O}(10^{-8} - 10^{-10})$, while the predicted range for non-perturbative long-distance contributions\(^3\) is approximately bounded by the box diagram rate and the current experimental sensitivity of $\mathcal{O}(10^{-4})$. New physics predictions span the same large range.\(^4\) We present evidence for $D$ mixing consistent with these expectations and with previous experimental limits.\(^5\)\(^6\)\(^7\)\(^8\)\(^9\)\(^10\)\(^11\) We also compare $D^0$ and $\bar{D}^0$ samples separately, and find no evidence for $C\bar{P}$ violation.

The mixing rate is characterized using the right-sign (RS), Cabibbo-favored (CF) decay\(^a\) $D^0 \rightarrow K^-\pi^+$ and the wrong-sign (WS) decay $D^0 \rightarrow K^+\pi^-$. The WS final state can be produced either through a doubly Cabibbo-suppressed (DCS) tree-level decay or through mixing followed by a CF decay. The DCS decay has a small rate $R_D$ of order $\tan^4 \thetaC \approx 0.3\%$ relative to CF decay, where $\thetaC$ is the Cabibbo angle. We distinguish $D^0$ and $\bar{D}^0$ by their production in the decay $D^{*+} \rightarrow \pi^+_m D^0$. In RS decays, the $\pi^+_m$ and kaon have oppositely signed charges, while in WS decays the charge signs are the same. The time dependence of the WS decay rate is used to separate DCS from mixed decays.

Charm mixing is generally characterized by two dimensionless parameters, $x \equiv \Delta m/\Gamma$ and $y \equiv \Delta \Gamma/2\Gamma$, where $\Delta m = m_3 - m_1$ ($\Delta \Gamma = \Gamma_3 - \Gamma_1$) is the mass (width) difference between the two neutral $D$ mass eigenstates and $\Gamma$ is the average width. If either $x$ or $y$ is non-zero, then $D$ mixing will occur. We approximate the time dependence of the WS decay of a meson produced

\(^a\)The use of charge-conjugate modes is implied unless otherwise noted.
as a $D^0$ at time $t = 0$ in the limit of small mixing ($|x|, |y| \ll 1$) and CP conservation as

$$T_{\text{ss}} = R_D + \sqrt{2} R_D g' \Gamma t + \frac{x^2 + y^2}{4} (\Gamma t)^2,$$

where $x' = x \cos \delta_{K\pi} + y \sin \delta_{K\pi}$, $y' = -x \sin \delta_{K\pi} + y \cos \delta_{K\pi}$, and $\delta_{K\pi}$ is the strong phase between the DCS and CP amplitudes. We study both CP-conserving and CP-violating cases. For the CP-conserving case, we fit for the parameters $R_D$, $x''$, and $y''$. To search for CP violation, we apply Eq. 1 to the $D^0$ and $\bar{D}^0$ samples separately, fitting for the parameters $R_0^\pm$, $x'^2$, $y'^2$ for $D^0(\pm)$ and $\bar{D}^0(\pm)$ decays.

2 Event Selection and Analysis

We use 384 fb$^{-1}$ of $e^+e^-$ colliding-beam data recorded near $\sqrt{s} = 10.6$ GeV with the Babar detector$^{13}$ at the PEP-II asymmetric-energy storage rings. We initially select signal candidates by combining oppositely-charged tracks identified as $K$ or $\pi$ using a likelihood-based particle identification algorithm, requiring the $K^{+}\pi^{-}$ invariant mass $1.81 < m_{K\pi} < 1.92$ GeV/$c^2$ and $e^+e^-$ center-of-mass frame (CM) momentum $p_{\text{CM}} > 2.5$ GeV/c. We require the $\pi^0_\text{fs}$ to have laboratory momentum $p > 0.1$ GeV/c and CM momentum $p^* < 0.45$ GeV/c.

To obtain the proper decay time $t$ and its error $\sigma_t$ for each $D^0$ candidate, we refit the $D^0$ daughter tracks and the $\pi^+_\text{fs}$, constraining the $D^0$ daughters to originate from a common vertex while simultaneously requiring the $D^0$ and $\pi^+_\text{fs}$ to originate from a common vertex constrained by the position and size of the $e^+e^-$ interaction region. We require a refit $\chi^2$ probability $P(\chi^2) > 0.1\%$, $M(K^+\pi^-) - m_{K\pi}$ mass difference $0.14 < \Delta m < 0.16$ GeV/$c^2$, proper decay time $-2 < t < 4$ ps and proper decay time error $\sigma_t < 0.5$ ps. The nominal $D^0$ mean proper lifetime is $\sim 0.410$ ps$^{14}$ and the most probable value of $\sigma_t$ for signal events is 0.16 ps. If there are multiple signal candidates with overlapping tracks within an event, we retain only the candidate with the highest $P(\chi^2)$. After applying all selection criteria, we retain approximately 1,229,000 RS and 64,000 WS signal candidates. To avoid potential bias, the complete event selection and analysis procedures were finalized prior to examining the mixing results.

The mixing parameters are determined using an unbinned, extended maximum-likelihood fit to the RS and WS data samples over the four observables $m_{K\pi}$, $\Delta m$, $t$, and $\sigma_t$. The fit is performed in several stages. First, the shapes of the RS and WS signal and background probability density functions (PDFs) are determined from an initial 2-d fit to $\{m_{K\pi}, \Delta m\}$. These shapes are then fixed in subsequent fits. Next, the $D^0$ proper-time resolution function and lifetime are determined from a fit to the RS data using $\{m_{K\pi}, \Delta m\}$ to separate the signal and background components. Finally, the WS data is analyzed using three different fit models. The first model assumes both CP conservation and the absence of mixing, the second model allows mixing but no CP violation, and the third model allows mixing and CP violation.

The RS and WS $\{m_{K\pi}, \Delta m\}$ distributions are described by four components: signal, random $\pi^+_\text{fs}$, misreconstructed $D^0$ and combinatorial background. The signal component has a characteristic peak in both $m_{K\pi}$ and $\Delta m$. The random $\pi^+_\text{fs}$ component models reconstructed $D^0$ decays combined with a random slow pion and has the same shape in $m_{K\pi}$ as signal events, but does not peak in $\Delta m$. Misreconstructed $D^0$ events have one or more of the $D^0$ decay products either not reconstructed or reconstructed with the wrong particle hypothesis. They peak in $\Delta m$, but not in $m_{K\pi}$. For RS events, most of these are semileptonic $D^0$ decays. For WS events, the main contribution is RS $D^0 \rightarrow K^-\pi^+$ decays where the $K^-$ and the $\pi^+$ are misidentified as $\pi^-$ and $K^+$, respectively. Combinatorial background events comprise the remainder of events and do not exhibit any peaking structure in $m_{K\pi}$ or $\Delta m$.

We fit the RS and WS data samples simultaneously to determine the PDF parameters describing the signal and random $\pi^+_\text{fs}$ event class shapes shared between RS and WS datasets.
We find $1,141,500 \pm 1,200$ RS signal events and $4,030 \pm 90$ WS signal events. The dominant background component is the random $\pi_s^+$ background. Projections of the WS data and fit are shown in Fig. 1.

![Figure 1](image)

Figure 1: a) $m_{K\pi}$ for WS candidates with $0.1445 < \Delta m < 0.1465$ GeV/c$^2$, and b) $\Delta m$ for WS candidates with $1.843 < m_{K\pi} < 1.883$ GeV/c$^2$. The fitted PDFs are overlaid.

The measured proper-time distribution for the RS signal is described by an exponential function convolved with a resolution function whose parameters are determined by the fit to the data. The resolution function is the sum of three Gaussians with widths proportional to the estimated event-by-event proper-time uncertainty $\sigma_t$. The random $\pi_s^+$ background is described by the same proper-time distribution as signal events, since the slow pion has little weight in the vertex fit. The proper-time distribution of the combinatorial background is described by a sum of two Gaussians, one of which has a power-law tail to account for a small long-lived component. The combinatorial background and real $D^0$ decays have different $\sigma_t$ distributions, as determined from data using a background-subtraction technique based on the $\{m_{K\pi}, \Delta m\}$ fit.

The fit to the RS proper-time distribution is performed over all events in the full $m_{K\pi}$ and $\Delta m$ region. The PDFs for signal and background in $m_{K\pi}$ and $\Delta m$ are used in the proper-time fit with all parameters fixed to their previously determined values. The fitted $D^0$ lifetime is found to be consistent with the world-average lifetime.

The measured proper-time distribution for the WS signal is modeled by Eq. 1 convolved with the resolution function determined in the RS proper-time fit. The random $\pi_s^+$ and misreconstructed $D^0$ backgrounds are described by the RS signal proper-time distribution since they are real $D^0$ decays. The proper-time distribution for WS data is shown in Fig. 2. The fit results with and without mixing are shown as the overlaid curves. The fit with mixing provides a substantially better description of the data than the fit with no mixing. The significance of the mixing signal is evaluated based on the change in negative log likelihood with respect to the minimum. Figure 3 shows confidence-level (CL) contours calculated from the change in log likelihood ($-2\Delta \ln \mathcal{L}$) in two dimensions ($x^2$ and $y'$) with systematic uncertainties included. The likelihood maximum is at the unphysical value of $x'^2 = -2.2 \times 10^{-4}$ and $y' = 9.7 \times 10^{-3}$. The value of $-2\Delta \ln \mathcal{L}$ at the most likely point in the physically allowed region ($x'^2 = 0$ and $y' = 6.4 \times 10^{-3}$) is 0.7 units. The value of $-2\Delta \ln \mathcal{L}$ for no-mixing is 23.9 units. Including the systematic uncertainties, this corresponds to a significance equivalent to 3.9 standard deviations (1−CL = $1 \times 10^{-4}$) and thus constitutes evidence for mixing. The fitted values of the mixing
parameters and $R_D$, along with errors, are listed in Table 1. The correlation coefficient between the $x'^2$ and $y'$ parameters is $-0.94$.

![Figure 2: a) Projections of the proper-time distribution of combined $D^0$ and $\bar{D}^0$ WS candidates and fit result integrated over the signal region $1.833 < m_{K^+} < 1.883$ GeV/$c^2$ and $0.1445 < \Delta m < 0.1465$ GeV/$c^2$. The result of the fit allowing (not allowing) mixing but not CP violation is overlaid as a solid (dashed) curve. b) The points represent the difference between the data and the no-mixing fit. The solid curve shows the difference between fits with and without mixing.]

![Figure 3: The central value (point) and confidence-level (CL) contours for $1 - \text{CL} = 0.317$ (1$, \sigma$), $4.55 \times 10^{-2}$ (2$, \sigma$), $2.70 \times 10^{-2}$ (3$, \sigma$), $6.33 \times 10^{-5}$ (4$, \sigma$) and $5.73 \times 10^{-7}$ (5$, \sigma$), calculated from the change in the value of $-2 \ln \mathcal{L}$ compared with its value at the minimum. Systematic uncertainties are included. The no-mixing point is shown as a plus sign (+).]

Allowing for the possibility of CP violation, we calculate the values of $R_D = \sqrt{R_D^+ R_D^0}$ and $A_D = (R_D^+ - R_D)/\sqrt{(R_D^+ + R_D)}$ listed in Table 1, from the fitted $R_D^\pm$ values. The best fit points $(x'^2, y'^2)$ shown in Table 1 are more than three standard deviations away from the no-mixing hypothesis. The shapes of the $(x'^2, y'^2)$ CL contours are similar to those shown in Fig. 3. All cross-checks indicate that the close agreement between the separate $D^0$ and $\bar{D}^0$ fit results is coincidental.

As a cross-check of the mixing signal, we perform independent \{m_{K^+}, \Delta m\} fits with no shared parameters for intervals in proper time selected to have approximately equal numbers
Table 1: Results from the different fits. The first uncertainty listed is statistical and the second systematic.

<table>
<thead>
<tr>
<th>Fit type</th>
<th>Parameter</th>
<th>Fit Results ($/10^{-3}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>No CP violation or mixing</td>
<td>$R_D$</td>
<td>$3.53 \pm 0.08 \pm 0.04$</td>
</tr>
<tr>
<td></td>
<td>$x'^2$</td>
<td>$3.03 \pm 0.16 \pm 0.10$</td>
</tr>
<tr>
<td></td>
<td>$y'$</td>
<td>$-0.22 \pm 0.30 \pm 0.21$</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$9.7 \pm 4.4 \pm 3.1$</td>
</tr>
<tr>
<td>CP violation allowed</td>
<td>$R_D$</td>
<td>$3.03 \pm 0.16 \pm 0.10$</td>
</tr>
<tr>
<td></td>
<td>$A_D$</td>
<td>$-21 \pm 52 \pm 15$</td>
</tr>
<tr>
<td></td>
<td>$x'^2+$</td>
<td>$-0.24 \pm 0.43 \pm 0.30$</td>
</tr>
<tr>
<td></td>
<td>$y'^+$</td>
<td>$9.8 \pm 6.4 \pm 4.5$</td>
</tr>
<tr>
<td></td>
<td>$x'^2-$</td>
<td>$-0.20 \pm 0.41 \pm 0.29$</td>
</tr>
<tr>
<td></td>
<td>$y'^-$</td>
<td>$9.6 \pm 6.1 \pm 4.3$</td>
</tr>
</tbody>
</table>

of RS candidates. Figure 4 shows the resulting fitted WS branching fractions growing with increasing proper time. The slope of a linear fit to the data points is consistent with the measured mixing parameters and inconsistent with the no-mixing hypothesis.

![Figure 4: The WS branching fractions from independent $m_{K^\tau}$, $\Delta m$ fits to slices in measured proper time (points). The dashed line shows the expected wrong-sign rate as determined from the mixing fit shown in Fig. 2. The $\chi^2$ with respect to expectation from the mixing fit is 1.3; for the no-mixing hypothesis (a constant WS rate), the $\chi^2$ is 24.9.](image)

We validated the fitting procedure on simulated data samples with both MC samples with the full detector simulation and large parameterized MC samples. In all cases we found the fit to be unbiased. As a further cross-check, we performed a fit to the RS data proper-time distribution allowing for mixing in the signal component; the fitted values of the mixing parameters are consistent with no mixing. In addition we found the staged fitting approach to give the same solution and confidence regions as a simultaneous fit in which all parameters are allowed to vary.

In evaluating systematic uncertainties in $R_D$ and the mixing parameters we considered variations in the fit model and in the selection criteria. We also considered alternative forms of the $m_{K^\tau}$, $\delta m$, proper time, and $\sigma_\ell$ PDFs. We varied the $t$ and $\sigma_\ell$ requirements. In addition, we considered variations that keep or reject all $D^{\pm+}$ candidates sharing tracks with other candidates. For each source of systematic error, we compute the significance $s^2_i = 2 \left[ \ln L(x'^2, y') - \ln L(x'^2_i, y'_i) \right] / 2.3$, where $(x'^2, y')$ are the parameters obtained from the standard fit, $(x'^2_i, y'_i)$ the parameters from the fit including the $i^{th}$ systematic variation, and $L$ the likelihood of the standard fit. The factor 2.3 is the 68% confidence level for 2 degrees of free-
dom. To estimate the significance of our results in \((x'^2, y')\), we reduce \(-2\Delta \ln L\) by a factor of \(1 + \Sigma a_i^2 = 1.3\) to account for systematic errors. The largest contribution to this factor, 0.06, is due to uncertainty in modeling the long decay time component from other \(D\) decays in the signal region. The second largest component, 0.05, is due to the presence of a non-zero mean in the proper time signal resolution PDF. The mean value is determined in the RS proper time fit to be 3.6 fs and is due to small misalignments in the detector. The error of \(15 \times 10^{-3}\) on \(A_D\) is primarily due to uncertainties in modeling the differences between \(K^+\) and \(K^-\) absorption in the detector.

We have presented evidence for \(D^0 - \bar{D}^0\) mixing. Our result is inconsistent with the no-mixing hypothesis at a significance of 3.9 standard deviations. We measure \(y' = [9.7 \pm 4.4 \text{ (stat.)} \pm 3.1 \text{ (syst.)}] \times 10^{-3}\), while \(x'^2\) is consistent with zero. We find no evidence for \(CP\) violation and measure \(R_D\) to be \([0.303 \pm 0.016 \text{ (stat.)} \pm 0.010 \text{ (syst.)}]\). The result is consistent with SM estimates for mixing.

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We are grateful for the extraordinary contributions of our PEP-II colleagues in achieving the excellent luminosity and machine conditions that have made this work possible. The success of this project also relies critically on the expertise and dedication of the computing organizations that support Babar. The collaborating institutions wish to thank SLAC for its support and the kind hospitality extended to them. This work is supported by the US Department of Energy and National Science Foundation, the Natural Sciences and Engineering Research Council (Canada), the Commissariat à l’Energie Atomique and Institut National de Physique Nucléaire et de Physique des Particules (France), the Bundesministerium für Bildung und Forschung and Deutsche Forschungsgemeinschaft (Germany), the Istituto Nazionale di Fisica Nucleare (Italy), the Foundation for Fundamental Research on Matter (The Netherlands), the Research Council of Norway, the Ministry of Science and Technology of the Russian Federation, Ministerio de Educación y Ciencia (Spain), and the Science and Technology Facilities Council (United Kingdom). Individuals have received support from the Marie-Curie IEF program (European Union) and the A. P. Sloan Foundation.

References

CLEO-c RESULTS

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Running at various charm thresholds the CLEO-c detector has provided the world's largest sample of $\psi(3770)$ and $\psi(4170)$ decays. Using these data-sets we have reconstructed $D_{(s)}$ mesons decaying to hadronic final states and used them to tag events with a charm-anti-charm meson pair. We present the precision extraction of the $D$ and $D_s$ meson decay constants, measurements of CKM matrix elements ($V_{cs}$ and $V_{cd}$), and semi-leptonic form factors. In addition, we investigate a variety of rare semi-leptonic decays.

1 Introduction

Leptonic and semileptonic decays of $B$ and $D$ mesons provide one of the best experimental frameworks for measurement of various standard model parameters. Pure leptonic decays are the classic way to determine the decay constants of charged pseudoscalar mesons. Semileptonic decays of heavy-light mesons are used to determine CKM matrix elements. The accuracy of $|V_{ub}|$ is limited by large theoretical uncertainty. Recently, dramatic progress has been achieved in lattice QCD, for a wide variety of hadronic quantities. Precision measurements of semileptonic $D_{(s)}$ decays combined with leptonic measurements so that the CKM matrix elements drops out, constitute a direct and vital check of lattice QCD methods for heavy quark systems and also provide an important experimental input for calculations of the $B$ meson decay constants.

2 Program Overview and CLEO-c Data Sets

The CLEO-c physics program spans the precision measurements involving mesons containing charm quarks, probes aspects of flavor mixing in the charm sector, and enables precision probes at other experiments. For the past 3 years the Cornell $e^+e^-$ storage ring (CESR) has become the charm factory and has produced sizeable statistics at 3770 and 4170 center of mass energies.

CLEO-c detector has collected data at various charm threshold energies. As of March 2007, CLEO collaboration has collected about $315\,pb^{-1}$, $281\,pb^{-1}$, and $50\,pb^{-1}$ at 4170 MeV, 3770 MeV, and 3686 MeV energies, respectively. In exploratory scan for maximum $D_s\bar{D}_s$ yield we accumulated $\sim60\,pb^{-1}$ at center of mass energies ranging from 3.97 to 4.26 GeV. Details of the CLEO-c program may be found elsewhere. Results presented here use about half of the luminosity we intend to acquire. By March 2008 CLEO collaboration will have the full data sample.

In all results quoted in this paper, the first uncertainty is due to statistics, the second uncertainty is due to systematics and the third uncertainty (if applicable) is due to theoretical
3 Physics at the $\psi(3770)$ Threshold

$\psi(3770)$ predominantly decays to $D\bar{D}$ mesons. CLEO's excellent geometrical acceptance enables us to fully reconstruct $D$ mesons in variety of hadronic modes. We analyze $D$-meson opposite the fully reconstructed “tagged” $D$. Tagged $D$ creates a beam of $D$ mesons with known momentum. From the remaining tracks and showers in the event a (semi)leptonic decay is reconstructed. The momentum of the $D$ tag is used as an estimate of the momentum of the $D$ decaying (semi)leptonically, which is not measured due to the undetected neutrino. Tagging can be used to find absolute branching fractions not only to hadronic decays, but semi-leptonic and even pure leptonic decays that can not be fully reconstructed.

$K_s$ and $\pi^0$ are reconstructed from displaced vertices and two detected photons in the CsI calorimeter, respectively. $dE/dx$ provides good $K-\pi$ separation up to 600 MeV/c. For higher momentum tracks we combine RICH information with $dE/dx$. For electron identification we require a match between the momentum measurement in the tracking system and the energy deposited in the CsI calorimeter as well.

With known initial 4-momentum of the resonance from the beam parameters, we impose energy and momentum conservation for the $D$ meson candidates. This is done by requiring the difference in the energy, $\Delta E = E_{\text{cand}} - E_{\text{beam}}$, is consistent with zero within 2.5 standard deviations and the beam-constrained mass of the $D$ meson candidate, $M_{bc} = \sqrt{E_{\text{beam}}^2 + p_{\text{cand}}^2}$ is consistent with nominal mass to within 3 standard deviations. The signal to noise ratio in such a distribution is excellent. We used the technique of “D-tagging”, as described above, in the current analysis. The net-tagging efficiency from a combination of $D$ modes is approximately 20%. We use a total of 160K single tag events for further analysis.

3.1 Leptonic Decays

The Standard Model decay rate for $D^+ \to l^+\nu$ is given by:

$$\Gamma(D^+_{(s)} \to l^+\nu) = \frac{G_F^2}{8\pi} f_{D^+_{(s)}} \frac{m_t}{m_{D^+_{(s)}}} (1 - \frac{m_l}{m_{D^+_{(s)}}})^2 |V_{cd(s)}|^2,$$  \hspace{1cm} (1)

where $M_{D^+_{(s)}}$ is the $D^+_{(s)}$ mass, $m_l$ is the mass of the final state lepton, $V_{cd}$ is a CKM matrix element, and $G_F$ is the Fermi coupling constant. The above relation holds for all pseudoscalar meson decays.

Branching Fraction: $D^+ \to \tau^+\nu_\tau$

This analysis follows very closely CLEO analysis for $D^+ \to \mu^+\nu^\tau$. We use the same sample of fully reconstructed charged $D$ decays for tags. Opposite this tag we search for a single charged track of opposite charge to the tag unaccompanied by additional charged tracks or significant neutral energy. Here we look for tracks consistent with $\pi^+$ candidates from the decay sequence $D^+ \to \tau^+\nu_\tau, \tau^+ \to \pi^+\nu$. The $MM^2$ distribution in Monte Carlo is shown in the left plot in Fig. 1. This search has a blurred signal region, as compared to muon case, because the small $D^+ - \tau^+$ mass difference causes the $\tau^+$ to be almost at rest in the lab frame and thus the $\pi^+$ has relatively large momentum. We also make sure that we do not accept $D^+ \to \mu^+\nu$ events or semi-leptonic decays with electrons.

In this analysis we consider three cases: (i) the track deposits < 300 MeV in the calorimeter, characteristic of a non-interacting pion or muon; (ii) track deposits > 300 MeV in the calorimeter,
characteristic of an interacting pion: (iii) the track is consistent with an electron. For all three cases we form a missing mass squared, $MM^2 = (E_{beam} - E_{track})^2 - (p_D^- - p_{track})^2$, where $E_{beam}$ is the beam energy, $E_{track}$ the measured track energy taking as the mass $\mu^+$, $\pi^+$, or $e^+$ for cases (i)-(iii), respectively, $p_{D^-}$ is the three-momentum of the fully reconstructed $D^-$. The track candidates are required to be within the barrel region of the detector. For cases (i) and (ii) we insist that the track is not consistent with kaon. For case (i) we define the signal region to be in the interval $0.175 > MM^2 > 0.05 \text{ GeV}^2$, we exclude the region close to zero $MM^2$, where $\mu^+\nu$ signal appears. While for case (ii) we define the signal region to be in the interval $0.175 > MM^2 > 0.05 \text{ GeV}^2$, here we specifically select pions therefore we extend the signal region. The $MM^2$ distributions for cases (i) and (ii) are shown in the left plot in Fig. 1(a) and (b). The peak at $MM^2 = 0.25$ is due to $D^+ \rightarrow \pi^+ K_L$ decays. We do not observe a statistically significant difference between signal and background events. For case (i) we have a net signal of $5.9 \pm 3.5 \pm 0.3 \tau^+\nu$ events. For case (ii) our yield is $3.0 \pm 2.9 \pm 0.2$ events. For both cases combined we have $8.9 \pm 4.6$ events, where the error includes both statistical and systematic components. Since the result is not statistically significant we quote an upper limit of $\mathcal{B}(D^+ \rightarrow \tau^+\nu) < 2.1 \times 10^{-3}$ at 90% confidence level. The ratio of $\Gamma(D^+ \rightarrow \tau^+\nu)/\Gamma(D^+ \rightarrow \mu^+\nu)$ to the expected Standard Model value using using CLEO-c measured $\mathcal{B}(D^+ \rightarrow \mu^+\nu)$ is $< 1.8$ at 90% confidence level.

3.2 Semi-Leptonic Decays

Improved measurements of inclusive lepton spectra from charm mesons are of considerable interest. In addition to the integrated spectra providing the inclusive branching fraction, the theoretical interpretation is cleaner as spectra are independent of the hadronic width, the detailed shape is useful when secondary charm leptons are a background to semileptonic B decay studies, and non perturbative effects are pronounced in the lepton endpoint region. We focus on electrons due to difficulties of soft muon identification at the lower CLEO-c energies.

Semileptonic decays of the type $D \rightarrow h\ell\bar{\nu}$ with a pseudoscalar hadron in the final state in the limit of small lepton masses is described as follows:

$$\frac{d\Gamma}{dq^2} \propto |f_h^A(q^2)|^2 \times p_{\ell}^2 \times |V_{eq}|^2,$$  (2)
where \( q^2 \) is the momentum transfer to the virtual \( W \), \( f_{s}^+(q^2) \) is the form factor function describing the probability to end up with a hadron of type \( h \) in the final state for a given \( q^2 \), \( p_h \) is the momentum of the outgoing hadron, \( V_{cq} \) is the corresponding CKM matrix element.

We employ two different methods described below, to extract exclusive branching fractions, form factors, and \( V_{cd} \) and \( V_{cs} \) matrix elements. The samples obtained using using these two methods have 40\% overlap.

**First Method “D-Tag”**

In this method, we look for the signal \( D \)-meson opposite the fully reconstructed “tagged” \( D \). We separate signal events from the background by defining \( U = E_{miss} - |\vec{p}_{miss}| \), where \( E_{miss} \) and \( \vec{p}_{miss} \) are the missing energy and momentum of the \( D \) decaying semileptonically, both of which are measured. If the massless particle is missing then the variable \( U \) peaks at zero. Events in which there are missing particles (with given mass) usually peak at large positive values of \( U \). Semileptonic decays in which there is an incorrect assignment of particles mostly populate the region outside the signal region. The left plot in Fig. 2 shows the \( U_{miss} \) distributions of \( D \rightarrow K^{(*)} e^+\nu \); Top left is \( D^0 \rightarrow K^- e^+\nu \); top right is \( D^+ \rightarrow K^0 e^+\nu \); bottom left is \( D^0 \rightarrow \pi^- e^+\nu \); and bottom right is \( D^+ \rightarrow \pi^0 e^+\nu \). Branching fractions are listed in Table 2 with a label “T” for the Tagged method.

**Second Method “Untagged \( \nu \)-Reconstruction”**

In this method, we form a candidate signal by combining a signal electron track and signal hadron (\( \pi^+ / K^+ \)) with the missing \( p_{miss}^\mu \) of the neutrino (\( = p_{beam}^\mu \) of beam minus \( p_{vis}^\mu \) of the other charged tracks and visible shower energies, where \( p^\mu \) is the 4-momentum). Finally, we form the beam constrained \( D \) mass in the center of mass rest frame, as, \( M_{bc} = \sqrt{E_{beam}^2 - (p_{K(\pi)} + p_e + \xi p_{miss})^2} \), where \( E_{beam}, p_{K(\pi)}, p_e, \) and \( p_{miss} \) are the beam energy, \( K(\pi) \) 3-momentum, electron 3-momentum, and neutrino 3-momentum, respectively. \( \xi \) is the scale factor such that \( \Delta E \equiv E_D - E_{beam} = 0 \). The beam constrained mass (\( M_{bc} \)) is shown in the right plot in Fig. 2: Top left is \( D^0 \rightarrow \pi^- e^+\nu \); top right is \( D^0 \rightarrow K^- e^+\nu \); bottom left is \( D^+ \rightarrow \pi^0 e^+\nu \); and bottom right is \( D^+ \rightarrow K^0 e^+\nu \). Branching fractions are listed in Table 2 with a label “U” for the Untagged method.

**Form Factors and CKM Matrix Element Extractions**

For each signal decay candidate we calculate the invariant mass squared of the \( e^+\nu \) pair (virtual \( W \), \( q^2 \)). We analyze our semileptonic sample in bins of \( q^2 \) using the two different methods discussed above. We perform two different fits to extract form factors and CKM matrix elements.

Theoretical predictions for the form factor shape can be tested against the \( q^2 \) distribution in data, corrected for the \( q^2 \)-dependent detection efficiency. The normalization is determined by the product \( [V_{cq} f_c(0)] \). Due to the precision with which \( V_{cs} \) and \( V_{cd} \) are determined in other experiments (2 – 4\%), a measurement of \( [V_{cq} f_c(0)] \) determines \( f_c(0) \). The simplest reasonable parameterization of the pseudoscalar form factors is referred as single pole shape or “vector-meson dominance”, \( , \sim 1/(1 - q^2/M_{pole}^2) \), where \( M_{pole} \) is the mass of the pole, which is expected to be \( D^0 \pi^+ (D^0 K^+) \) vector state. More sophisticated models are: the modified model which explicitly incorporates the \( D^0 \) mass but includes a term \( a \) to account for deviations from the vector masses; the Hill series parameterization is a less model-dependent way of dealing with the analytic singularities at \( q^2 = M_{pole}^2 \). Table 1 lists CLEO-c measured results and a comparison with previous CLEOII \( f_d \) and LQCD results. Form factor fits are shown in Fig. 3 and the fit results are shown in Table 2.
Figure 2: Left plot. Tag Method: Fit to $U$ for $D \to he^{+}\nu$ distributions. Points with error bars are data. Red is the fitted curve to signal and backgrounds, and blue, yellow and green are fits to various backgrounds. Right plot. Un-tagged method: Fit $M_{D^{0}}$ distributions for $D \to he^{+}\nu$. Points represent data and stacked histograms are: Clear - signal Monte Carlo, Gray - summed background Monte Carlo, Light Gray - fakes from data.

Table 1: Simple pole and modified pole preliminary fit results from $D \to K(x)e^{+}\nu$ decays. Tag and untag results are not to be average.

<table>
<thead>
<tr>
<th></th>
<th>$M_{pole}^{D^+\rightarrow Ke^{+}}$</th>
<th>$M_{pole}^{D^+\rightarrow \pi e^{+}}$</th>
<th>$\alpha_{D^+\rightarrow Ke^{+}}$</th>
<th>$\alpha_{D^+\rightarrow \pi e^{+}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Tag</td>
<td>1.96 ± 0.03 ± 0.01</td>
<td>1.95 ± 0.04 ± 0.02</td>
<td>0.22 ± 0.05 ± 0.02</td>
<td>0.17 ± 0.10 ± 0.05</td>
</tr>
<tr>
<td>Untagged</td>
<td>1.97 ± 0.03 ± 0.01</td>
<td>1.89 ± 0.04 ± 0.01</td>
<td>0.21 ± 0.04 ± 0.03</td>
<td>0.32 ± 0.07 ± 0.03</td>
</tr>
<tr>
<td>CLEO III</td>
<td>1.89 ± 0.05 ± 0.04</td>
<td>1.86 ± 0.06 ± 0.03</td>
<td>0.36 ± 0.10 ± 0.03</td>
<td>0.37 ± 0.07 ± 0.03</td>
</tr>
<tr>
<td>LQCD</td>
<td>1.72 ± 0.18 ±</td>
<td>1.99 ± 0.17</td>
<td>0.50 ± 0.04</td>
<td>0.44 ± 0.04</td>
</tr>
</tbody>
</table>

Using the unquenched lattice QCD result \(^3\) for $|f_{+}(0)|$ allows for an experimental measurement of the CKM matrix elements $V_{cd}$ and $V_{cs}$. The measured $V_{cd}$ and $V_{cs}$ for tagged (untagged) method yields the values $0.234 \pm 0.01 \pm 0.004 \pm 0.024 (0.229 \pm 0.007 \pm 0.005 \pm 0.024)$ and $1.014 \pm 0.013 \pm 0.009 \pm 0.106 (0.996 \pm 0.008 \pm 0.015 \pm 0.104)$, respectively. Our measurements are in good agreement with current world averages. $V_{cd}$ and $V_{cs}$ are the most precise and robust measurements using semileptonic decays to date. However, the uncertainties ($\sim 10\%$ relative) on the form factor magnitude dominate.

3.3 Rare Modes

CLEO-c has also looked at rare semileptonic decay modes of $D$-mesons. Table 3 tabulates the preliminary branching fractions. With the exception of $B(D^{+} \to \omega e^{+}\nu)$, these results are either first observation or improved upper limits. These results are consistent with predictions based on ISGW2 \(^{17}\) and heavy quark symmetry (HQSS) \(^{18}\).

4 Physics at the 4170 MeV Threshold

The uare leptonic decay widths for $D_{s}$ system are also defined by eq. 1, replacing the $d$-quark by an $s$-quark. Here we present the most precise measurement to date of $f_{D_{s}^{+}}$ \(^{9}\) and the ratio $f_{D_{s}^{+}}/f_{D^{+}}$. 
Figure 3: A comparison of LQCD and CLEO’s tagged data fits: Dashed fits are the LQCD and dots with error bars are the data points and solid lines represent fit to data. The left plot is \( D^0 \rightarrow \pi^- \tau^+ \nu \) and the right plot is \( D^0 \rightarrow K^- \nu \). 

Table 2: Fit Results. T implies tagged analysis and U means untagged \( \nu \)-reconstruction method. Form factors are derived using the 2006 global fit values for \( V_{td} \) and \( V_{cd} \). Tagged and untagged results are preliminary and are not to be averaged.

| Mode          | \( Br \% \) Tagged | \( \frac{|f_{2\pi}^{(K)}(0)|}{f_{2\pi}^{(\pi)}(0)} \) | CLEO-c | LQCD |
|---------------|---------------------|---------------------------------|--------|------|
| \( D^0 \rightarrow \pi^- e^+ \nu \) | T: 0.309 ± 0.012 ± 0.006 | U: 0.301 ± 0.011 ± 0.010 | T: 0.660 ± 0.028 ± 0.011 | \( 0.64 \pm 0.03 \pm 0.06 \) |
| \( D^0 \rightarrow \pi^0 e^+ \nu \) | T: 0.397 ± 0.027 ± 0.028 | U: 0.383 ± 0.025 ± 0.016 | U: 0.636 ± 0.017 ± 0.013 | \( 0.761 \pm 0.010 \pm 0.007 \) |
| \( D^0 \rightarrow K^- e^+ \nu \) | T: 3.58 ± 0.05 ± 0.05 | U: 3.56 ± 0.03 ± 0.11 | T: 0.749 ± 0.005 ± 0.010 | \( 0.73 \pm 0.03 \pm 0.07 \) |
| \( D^+ \rightarrow K^+ e^+ \nu \) | T: 8.86 ± 0.17 ± 0.20 | U: 8.75 ± 0.13 ± 0.30 | U: 0.636 ± 0.017 ± 0.013 | \( 0.64 \pm 0.03 \pm 0.06 \) |

We analyze both \( D_s^+ \rightarrow \mu^+ \nu \) and \( D_s^+ \rightarrow \tau^+ \nu \). Both \( D_s \) decays are helicity suppressed because the \( D_s \) is a spin-0 particle, and the final state consists of a naturally left-handed spin-1/2 neutrino and naturally right-handed spin-1/2 anti-lepton.

At \( \psi(4170) \) threshold the cross-section for \( D_s^+ D_s^- + D_s^+ D_s^- \) is \( \sim 1 \) nb, with \( D_s^+ D_s^- \) production being only \( \sim 5\% \) of this rate. \( D \) mesons are also produced mostly as \( D^* \bar{D}^* \), with a cross-section of \( \sim 5 \) nb. There is also a comparable contribution coming from \( D D^* + \bar{D} \bar{D}^* \) and \( D \bar{D} \) with extra pions. The underlying light quark “continuum” background is about \( 12 \) nb. The relatively large cross-sections, relatively large branching fractions and sufficient luminosities, allow us to fully reconstruct one \( D_s \) as a “tag” and examine the leptonic decays of the other \( D_s \) in the events of type \( e^+ e^- \rightarrow D_s^+ D_s^- \) or \( D_s^+ D_s^- \). We reconstruct tags from both final states.

4.1 Determination of Branching fractions and \( f_{D_s} \) from \( D_s^+ \rightarrow \mu \nu \) and \( \tau \nu \)

We employ two separate methods to extract branching fractions and decay constants. We find 19185±326 tagged events in our \( D_s \) sample and are used for further analysis.
First Method

As above, the missing mass squared is examined to identify the signal. The lone signal side track is required to not be an electron but instead to be muon or pion. An energy requirement of 300 MeV accepts 98% of all muon tracks and 60% of all pion tracks. We distinguish cases with an energy deposition below 300 MeV in the calorimeter (typical for $D_s^+ \rightarrow \mu^+ \nu_\mu$, but possible for $\tau \rightarrow \pi \nu_\tau$) from those that have above 300 MeV (in which case the track is likely to be a pion from $\tau \rightarrow \pi \nu_\tau$). The data show a clear enhancement in the expected $D_s^+ \rightarrow \mu^+ \nu_\mu$ signal region around $MM^2 = 0$, as shown in Fig. 4(i) (left). The $D_s^+ \rightarrow \tau \nu_\tau$ events are more spread out due to the presence of the additional, unreconstructed, neutrino, but signal events are seen as well, as shown in Fig. 4(ii) (left). The signatures for $\ell = \mu$ and $\tau$ overlap, hence the summed distribution is fit for the two branching fractions, where the SM ratio for the two is assumed. CLEO’s preliminary result, combining the two, is $B(D_s^+ \rightarrow \mu^+ \nu_\mu) = (0.664 \pm 0.076 \pm 0.028)$, or $f_{D_s} = (282 \pm 16 \pm 7)$ MeV. If the track is required to be consistent with an electron, no candidates are found, resulting in an upper limit of $B(D_s^+ \rightarrow e^+ \nu_e) < 3.1 \times 10^{-4}$.

Second Method

The second approach, based on the same data, uses the decay chain $\tau \rightarrow e \bar{\nu}_e \nu_\tau$. Its product with the $D_s^+ \rightarrow \tau^+ \nu_\tau$ branching fraction ($\sim 6 - 7\%$) is about 1.3%, to be compared with the inclusive semileptonic branching ratio $D_s^+ \rightarrow X e^+ \nu_e \sim 8\%$. The analysis procedure demands a sole electron-like track on the signal side, and limits the energy not associated with the other identified decay products in the calorimeter to be less than 400 MeV. No additional energy deposition is expected for signal events other than the transition photon from $D_s^+ D_s \rightarrow (\gamma D_s) D_s$, upon which no selection requirements are placed, and showers resulting from interactions of the decay products of the tag side with the detector material. The signal distribution, together with background estimates from Monte Carlo simulations, is presented in Fig. 4 (right). This analysis leads to $B(D_s^+ \rightarrow \tau^+ \nu_\tau) = (6.3 \pm 0.8 \pm 0.5\%)$ and $f_{D_s} = (278 \pm 17 \pm 12)$ MeV (both preliminary). Since the two measurements are complementary, one can form the average $f_{D_s} = (280 \pm 12 \pm 6)$ MeV and also use them to measure the ratio $B(D_s^+ \rightarrow \tau^+ \nu_\tau) / B(D_s^+ \rightarrow \mu^+ \nu_\mu) = (9.9 \pm 1.9)$, consistent with the SM expectation of 9.72.

5 Conclusions

CLEO-c results for leptonic and semileptonic decays of charm mesons, with already improved accuracy, provide stringent tests of theoretical model predictions and provide a good calibration for LQCD to apply their models in the beauty meson sector. More rare $D$-decay modes are accessible with large data sample. By March 2008, CLEO will further improve the precision with more data.

Table 3: Preliminary CLEO branching fraction measurements of rare semileptonic $D$ decays.

<table>
<thead>
<tr>
<th>Decay</th>
<th>CLEO result ($10^{-4}$)</th>
<th>PDG ($10^{-4}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^+ \rightarrow \omega t^+ \nu$</td>
<td>$14.9 \pm 2.7 \pm 0.5$</td>
<td>$16^{+7}_{-5}$</td>
</tr>
<tr>
<td>$D^+ \rightarrow \phi t^+ \nu$</td>
<td>$&lt; 2$ (90% CL)</td>
<td>$&lt; 209$</td>
</tr>
<tr>
<td>$D^+ \rightarrow \eta t^+ \nu$</td>
<td>$12.9 \pm 1.9 \pm 0.7$</td>
<td>$&lt; 70$</td>
</tr>
<tr>
<td>$D^+ \rightarrow \eta' t^+ \nu$</td>
<td>$&lt; 3$ (90% CL)</td>
<td>$&lt; 110$</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^- \pi^+ \pi^- t^+ \nu^8$</td>
<td>$2.9^{+1.0}_{-1.0} \pm 0.5$</td>
<td>$&lt; 12$</td>
</tr>
<tr>
<td>$K_{1}(1270) t^+ \nu$</td>
<td>$2.2^{+1.1}_{-1.0} \pm 0.2$</td>
<td>$-$</td>
</tr>
</tbody>
</table>
Figure 4: CLEO data on leptonic $D_0$ decays (preliminary). Left: Missing-mass-squared distributions from data corresponding to (i) $D^+_s \rightarrow \mu^+\nu_\mu + \tau^+\nu_\tau$, (ii) $D^+_s \rightarrow \tau^+\nu_\tau$, (iii) $D^+_s \rightarrow e^+\nu_e$. Right: $D^+_s \rightarrow \tau^+\nu_\tau$: Energy deposited in the calorimeter for tagged $D_s$ events and a single electron-like track on the signal side not accounted for by the tag or the electron candidate. The circles are data; the curves are MC predictions for signal as well as several semileptonic background sources.

6 Acknowledgments

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References

Recent BES results and future prospects

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Recent results on light hadron spectroscopy are reviewed, which include the observation of $X(1810)$ in $\omega \phi$ mass spectrum in $J/\psi \rightarrow \gamma \omega \phi$, the observation of a broad $1^{--}$ resonance of $K^+ K^-$ mass in $J/\psi \rightarrow K^+ K^- \pi^0$ and the production of $\sigma$ and $\kappa$. The first measurement of $\psi(2S)$ radiative decays, as well as the measurement of $\psi(3770)$ non-$D \bar{D}$ decays are presented too. The data samples used for these analyses consist of $5.8 \times 10^7 J/\psi$ events, $1.4 \times 10^7 \psi(2S)$ events and 33 pb$^{-1}$ data at $\psi(3770)$. We also report the status of BESIII/BEPCII project.

1 Introduction

One of the distinctive features of QCD as a non-Abelian gauge theory is the self-interaction of gluons, which predicts the existence of the multi-quark states, $q\bar{q}$-gluon hybrids and glueballs. These states have been searched for many years by some experiments. However, none of them is identified after all the efforts. The $5.8 \times 10^7 J/\psi$ events, accumulated with BES detector, provide a good laboratory for the search for non-$q\bar{q}$ states and study of light hadron spectroscopy. Here, we report the observation of $X(1810)$ in $\omega \phi$ mass spectrum in $J/\psi \rightarrow \gamma \omega \phi$, the observation of a broad $1^{--}$ resonance of $K^+ K^-$ mass in $J/\psi \rightarrow K^+ K^- \pi^0$ and the production of $\sigma$ and $\kappa$.

The radiative decays of $\psi(2S)$ have been limited due to low statistics in previous experiments. The measured branching ratios for radiative decays only sum up to about 0.05%, which is much lower than what expected. Based on $1.4 \times 10^7 \psi(2S)$ events, the first measurements of some $\psi(2S)$ radiative decays are presented.

The $\psi(3770)$ was considered to decay almost entirely to pure $D \bar{D}$ because its mass is above the open charm-pair threshold and its width is two orders of magnitude larger than that of $\psi(2S)$. Since the first observation of $\psi(3770)$ non-$D \bar{D}$ decays to $\pi \pi J/\psi$ by BES, many searches for exclusive $\psi(3770)$ non-$D \bar{D}$ decays have been performed by CLEO and BES. We present the measurement of $\psi(3770)$ non-$D \bar{D}$ decay branching ratio from BESII 33 pb$^{-1}$ $\psi(3770)$ data.
2 Light hadron spectroscopy

2.1 The $\omega\phi$ threshold enhancement $X(1810)$ in $J/\psi \rightarrow \gamma\omega\phi$

An enhancement near the $\omega\phi$ mass threshold is observed in the OZI suppressed decays of $J/\psi \rightarrow \gamma\omega\phi$, based on a sample of $5.8 \times 10^7 J/\psi$ events collected with the BESII detector. Figure 2 shows the $\omega\phi$ invariant mass distribution and a structure peaked near $\omega\phi$ threshold is observed. The dashed curve in the figure indicates how the acceptance varies with invariant mass.

![Figure 1: The \(\omega\phi\) invariant mass distribution. The dashed curve shows the acceptance curve.](image)

A partial wave analysis shows that the $J^P$ of this enhancement favors $0^+$ and its mass and width are $M = 1812^{+10}_{-8}(\text{stat})\pm 18(\text{syst}) \text{MeV}/c^2$ and $\Gamma = 105 \pm 20 (\text{stat}) \pm 28 (\text{syst}) \text{MeV}/c^2$. The product branching fraction is determined to be $B(J/\psi \rightarrow \gamma X) \cdot B(X \rightarrow \omega\phi) = (2.61 \pm 0.27(\text{stat}) \pm 0.65(\text{syst})) \times 10^{-4}$. The mass and width of this state are not compatible with any known scalars listed in PDG. However, more statistics and further studies are needed to clarify this enhancement.

2.2 Observation of a broad $1^{--}$ resonance in $J/\psi \rightarrow K^+K^-\pi^0$

The Dalitz plot for the selected $J/\psi \rightarrow K^+K^-\pi^0$ events is shown in Figure 2(a), where a broad $K^+K^-$ band is evident in addition to the $K^{*+}(892)$ and $K^{*+}(1410)$ signals. This band corresponds to the broad peak observed around 1.5 GeV/c$^2$ in the $K^+K^-$ invariant mass projection shown in Figure 2(b).

![Figure 2: (a) The Dalitz plot for $KK\pi^0$ candidate events. (b) The $KK$ invariant mass distribution for $KK\pi^0$ candidate events; the solid histogram is data and the shaded histogram is the background (normalized to data).](image)

A partial wave analysis shows that the $J^{PC}$ of this structure is $1^{--}$. Its pole position is determined to be $(1576^{+49}_{-55}(\text{stat})^{+98}_{-91}(\text{syst})) \text{MeV}/c^2 - i(409^{+111}_{-12}(\text{stat})^{+32}_{-67}(\text{syst})) \text{MeV}/c^2$. These parameters are not compatible with any known meson resonances.
2.3 The $\sigma$ and $\kappa$ production

The existence of $\sigma$ and $\kappa$ has been controversial. BES studied $J/\psi \to \omega \pi^+\pi^-$ decays\(^7\) based on 58M $J/\psi$ events. Figure ?? shows the $\pi^+\pi^-$ invariant mass spectrum recoiling against $\omega$. In addition to the well known $f_2(1270)$, a broad bump in the low mass region is clearly seen. Two independent partial wave analyses (PWA) are performed and different parametrizations of $\sigma$ amplitude are used. All give the consistent results for $\sigma$ pole position. The averaged pole is determined to be $(541 \pm 39 - i (252 \pm 42))$ MeV/c\(^2\).

Based on 14M $\psi(2S)$ events collected at BESII, a partial wave analysis is performed to $\psi(2S) \to \pi^+\pi^- J/\psi$, with $J/\psi \to \mu^+\mu^-$.\(^7\) A severe suppression of the $\pi^+\pi^-$ invariant mass near the $\pi\pi$ threshold is distinctively different from the phase space shape, which suggests the $\sigma$ production. Using different parametrizations of $\sigma$ amplitude, the data can be fitted well through a strong cancellation between $\sigma$ and a contact term. The obtained pole position is $(554\pm14\pm53) - i (242 \pm 5 \pm 24))$ MeV/c\(^2\), which is consistent with that from $J/\psi \to \omega \pi^+\pi^-$.\(^7\)

![Figure 3: The $\pi^+\pi^-$ invariant mass in $J/\psi \to \omega \pi^+\pi^-$.](image)

The $\kappa$ is studied in $J/\psi \to K^* K \pi$ with partial wave analysis at BES\(^7\). Fig. ?? is the $K\pi$ invariant mass spectrum which recoils against $K^*(892)$. The crosses are data and histograms represent the PWA fit projection. The shaded area shows the $\kappa$ contribution. The pole of $\kappa$ from $J/\psi \to K^* K \pi$ is determined to be $(841 \pm 30 \pm 81 / c^2) - i (309 \pm 45 \pm 48) \text{MeV/c}^2$.

![Figure 4: The $K\pi$ invariant mass recoiling against $K^*$. The crosses are data and histograms represent the PWA fit projection. The shaded area shows the $\kappa$ contribution.](image)

3 First measurement of $\psi(2S)$ radiative decays

The first measurements of branching fractions or upper limits for $\psi(2S)$ decays into $\gamma p\bar{p}$, $\gamma^2 (\pi^+\pi^-)$, $\gamma K_s K \pi$, $\gamma K^+ K^- \pi^+ \pi^-$, $\gamma K^{*0} K^- \pi^+ + c.c.$, $\gamma K^{*0} K^{*0}$, $\gamma \pi^+ \pi^- p\bar{p}$, $\gamma 2(K^+ K^-)$, $\gamma 3(\pi^+ \pi^-)$, and $\gamma 2(\pi^+ \pi^-) K^+ K^-$ with the invariant mass of hadrons below 2.9 GeV/c\(^2\) are performed\(^7\). The branching fractions of $\psi(2S)$ decays into $2(\pi^+ \pi^-) \pi^0$, $\omega \pi^+ \pi^-$, $\omega f_2(1270)$, $b^+_1 \pi^+$, and $\pi^{00}(\pi^+ \pi^-) K^+ K^-$ are also measured. Table ?? summarizes the branching fractions or upper limits for the $\psi(2S)$ radiative decays analyzed.
Table 1: Results for $\psi(2S) \rightarrow \gamma + \text{hadrons}$. For each final state, the following quantities are given: the number of events for $m_{X} < 2.9 \text{GeV}/c^{2}$ in $\psi(2S)$ data, $N^{X\text{ex}}$; the number of background events from $\psi(2S)$ decays and QED processes, $N^{\text{BG}}$; the number of signal events, $N^{\text{SG}}$; and the weighted averaged efficiency, $\epsilon$; the branching fraction with statistical and systematic errors or the upper limit on the branching fraction at the 90% C.L. Possible interference effects for the modes with intermediate states are ignored.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$N^{X\text{ex}}$</th>
<th>$N^{\text{BG}}$</th>
<th>$N^{\text{SG}}$</th>
<th>$\epsilon(%)$</th>
<th>$B(\times 10^{-5})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\gamma\rho\rho$</td>
<td>329</td>
<td>187</td>
<td>142 ± 18</td>
<td>35.3</td>
<td>2.9 ± 0.4 ± 0.4</td>
</tr>
<tr>
<td>$\gamma 2(\pi^{+}\pi^{-})$</td>
<td>1697</td>
<td>1114</td>
<td>583 ± 41</td>
<td>10.4</td>
<td>39.6 ± 2.8 ± 5.0</td>
</tr>
<tr>
<td>$\gamma K_{S}^{0}K^{+}K^{-} + c.c.$</td>
<td>$-$</td>
<td>$-$</td>
<td>115 ± 16</td>
<td>4.83</td>
<td>25.6 ± 3.6 ± 3.6</td>
</tr>
<tr>
<td>$\gamma K^{+}K^{-} + c.c.$</td>
<td>361</td>
<td>229</td>
<td>132 ± 19</td>
<td>4.94</td>
<td>19.1 ± 2.7 ± 4.3</td>
</tr>
<tr>
<td>$\gamma K^{0}K^{+}K^{-} + c.c.$</td>
<td>$-$</td>
<td>$-$</td>
<td>237 ± 39</td>
<td>6.86</td>
<td>37.0 ± 6.1 ± 7.2</td>
</tr>
<tr>
<td>$\gamma K^{+}K^{0}$</td>
<td>58</td>
<td>17</td>
<td>41 ± 8</td>
<td>2.75</td>
<td>24.0 ± 4.5 ± 5.0</td>
</tr>
<tr>
<td>$\gamma\pi^{+}\pi^{-}\rho\rho$</td>
<td>55</td>
<td>38</td>
<td>17 ± 7</td>
<td>4.47</td>
<td>2.8 ± 1.2 ± 0.5</td>
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<tr>
<td>$\gamma K^{+}K^{-}K^{0}$</td>
<td>15</td>
<td>8</td>
<td>$&lt; 14$</td>
<td>2.93</td>
<td>$&lt; 4.0$</td>
</tr>
<tr>
<td>$\gamma 3(\pi^{+}\pi^{-})$</td>
<td>118</td>
<td>95</td>
<td>$&lt; 45$</td>
<td>1.97</td>
<td>$&lt; 17$</td>
</tr>
<tr>
<td>$\gamma 2(\pi^{+}\pi^{-})K^{+}K^{-}$</td>
<td>17</td>
<td>13</td>
<td>$&lt; 15.5$</td>
<td>0.69</td>
<td>$&lt; 22$</td>
</tr>
</tbody>
</table>

4 $\psi(3770)$ non-$D\bar{D}$ decays

With the data taken at BESII, the $R$ values for inclusive hadronic event production at the center-of-mass energies of 3.650, 3.665 and 3.773 GeV are measured. The lowest order cross section for $\psi(3770)$ production is determined to be $\sigma_{\psi(3770)}^{J} = (9.575 \pm 0.256 \pm 0.813) \text{nb}$ at 3.773 GeV, the branching fractions for $\psi(3770)$ decays are $BF(\psi(3770) \rightarrow D^{0}\bar{D}^{0}) = (48.9 \pm 1.2 \pm 3.8)\%$, $BF(\psi(3770) \rightarrow D^{+}D^{-} = (35.0 \pm 1.1 \pm 3.3)\%$ and $BF(\psi(3770) \rightarrow D\bar{D}) = (83.9 \pm 1.6 \pm 5.7)\%$, which result in the total non-$D\bar{D}$ branching fraction of $\psi(3770)$ decay to be $BF(\psi(3770) \rightarrow non - D\bar{D}) = (16.1 \pm 1.6 \pm 5.7)\%$.

By analyzing the line-shapes of the cross sections for inclusive hadron, $D^{0}\bar{D}^{0}$ and $D^{+}D^{-}$ event production in the range from 3.660 GeV to 3.872 GeV covering both $\psi(2S)$ and $\psi(3770)$ resonances, we extract the branching fractions for $\psi(3770)$ decay into $D^{0}\bar{D}^{0}$ and $D^{+}D^{-}$ respectively to be $B(\psi(3770) \rightarrow D^{0}\bar{D}^{0}) = (46.7 \pm 4.7 \pm 2.3)\%$ and $B(\psi(3770) \rightarrow D^{+}D^{-} = (36.9 \pm 3.7 \pm 2.8)\%$, which give $B(\psi(3770) \rightarrow D\bar{D}) = (83.6 \pm 7.3 \pm 4.2)\%$ and non-$D\bar{D}$ branching fraction of $\psi(3770)$ to be $B(\psi(3770) \rightarrow non - D\bar{D}) = (16.4 \pm 7.3 \pm 4.2)\%$.

5 Search for $\eta'/\eta'$ invisible decays from $J/\psi \rightarrow \phi\eta/\eta'$

Using a data sample of $5.8 \times 10^{7}$ $J/\psi$ decays collected with the BES II detector at the BEPC, searches for invisible decays of $\eta$ and $\eta'$ in $J/\psi$ to $\phi\eta$ and $\phi\eta'$ are performed. No signals are found for the invisible decays of either $\eta$ or $\eta'$, and upper limits at the 90% confidence level are determined to be $1.65 \times 10^{-3}$ for the ratio $B(\eta \rightarrow invisible)/B(\eta \rightarrow \gamma\gamma)$ and $6.69 \times 10^{-2}$ for $B(\eta' \rightarrow invisible)/B(\eta' \rightarrow \gamma\gamma)$. These are the first searches for $\eta$ and $\eta'$ decays into invisible final states.

6 BESIII/BEPCII project

Beijing Electron Positron Collider (BEPC) and the detector BEijing Spectrometer (BES) were constructed and started to take data in 1989. At the end of 2003, the proposal of upgrading BEPC/BESII was approved by Chinese government and the construction was started. The accelerator has two storage rings with a circumference of 224 m, one for electron and one for positron, each with 93 bunches spaced by 8 ns. The peak luminosity is expected to be $10^{33} \text{cm}^{-2}\text{s}^{-1}$ at the
beam energy of 1.89 GeV. At the moment, the LINAC has been installed and tested successfully. The storage rings have been installed and is commissioning for synchrotron radiation run.

BESIII detector consists of a drift chamber (MDC) which has a small cell structure filled with a helium-based gas, an electromagnetic calorimeter (EMC) made of CsI(Tl) crystals, time-of-flight counters (TOF) for particle identification made of plastic scintillators, a muon system made of Resistive Plate Chambers (RPC) and a super conducting magnet. The designed single wire spatial resolution, $dE/dx$ resolution and momentum resolution for the MDC are 130 $\mu$m, 6% and 0.5% $\sqrt{1+p^2}$ ($p$ in GeV/c), respectively. The time resolution for TOF is $\sigma_{TOF} = 100$ ps for Bhabha events and the energy resolution for photons is $\sigma_E/E \simeq 2.5\% / \sqrt{E}$ (E in GeV) in EMC. The mass production of detector components is underway. Some have already been installed and are being tested. We expect to complete the detector installation at the end of this year and to start the data taking in 2008.

Acknowledgments

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Beyond SM, Extradim

K Physics
RECENT RESULT OF THE NA48 EXPERIMENT

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Recent results from different phases of the NA48 experiment will be reported. From the 2003-04 data taking NA48/2 the measurements of the CP asymmetry in $K^+ \rightarrow \pi^+\pi^+\pi^-$ and $K^+ \rightarrow \pi^+\pi^0\pi^0$, the first measurement of direct emission and interference term in $K^+ \rightarrow \pi^+\pi^0\gamma$ and the first observation of $K^+ \rightarrow \pi^+e^+e^-$ will be described. In addition the extraction of $V_{ub}$ from semileptonic decays, and a test of lepton universality using $K^+ \rightarrow e^+\nu$ and $K^+ \rightarrow \mu^+\nu$ will be addressed. Concerning NA48 and NA48/1 a measurement of the $K_L \rightarrow \pi^+\mu^+\nu$ form factor slope and the BR of $\Xi^0 \rightarrow \Lambda e^+e^-$ will be described.

1 Introduction

During 10 years of data taking NA48 has explored many topics in the charged and neutral Kaon physics, and continues to provide new results in both fields. In this paper will be discussed some of the most recent measurements produced by all three stages of the experimental program: NA48, NA48/1 and NA48/2. NA48 was devoted to the measurements of the CP violation in the neutral Kaon system through the parameter $\epsilon'/\epsilon$. The measurement has been performed during 3 years of data taking between 1997 and 2000 together with many studies of rare decays of both $K_L$ and $K_S$. The next stage NA48/1 in 2002 has been aimed to measure very rare $K_S$ decays $K_S \rightarrow \pi^0 e^+e^-$, $K_S \rightarrow \pi^0 \mu^+\mu^-$, and has produced also results in Hyperon physics. Finally the last stage NA48/2 has devoted his activity to the study of the CP violation in the charged Kaon sector. Together with the asymmetry measurements many other results in semileptonic and rare decays have been achieved.

2 Detector and beam line

The NA48 detector, used in all the described results, has been designed to measure with high precision momenta of both charged and neutral particles. The charged particle reconstruction is
provided by a magnetic spectrometer, with 4 drift chambers and a magnet, with a momentum resolution of \(\sigma(P)/P = (1.02 \pm 0.044\%)\), where \(P\) is in GeV/c; a charged hodoscope, with good time resolution, sends fast trigger signals representing the number of charged particles. The reconstruction of photon energy, direction, time, and position is given by a the LKr calorimeter. The calorimeter has an active volume of 10 m³ and an energy resolution of \(\sigma(E)/E = 3.2%/\sqrt{E} \pm 9%/E \pm 0.2\%\). Moreover it provides the position of the impinging photons with a precision of \((0.42/\sqrt{E}) \pm 0.06)\) cm. To identify muons from pions an hadronic calorimeter muon counter have been also constructed. Details on the detectors can be found in1.

The measurements presented in this paper have been obtained using different beam line configurations. The NA48 beam line was designed to produce and transport both \(K_L\) and \(K_S\) beams simultaneously1. The \(K_L\) beam was produced by SPS 450 GeV/c proton beam impinging on a beryllium target. The beginning of the decay volume was defined by the last of three collimators, located 126 m downstream of the target. The measurement \(K_L \to \pi^\pm\mu^\mp\nu\) form factors uses this beam configuration. For the NA48/1 experiment the \(K_L\) beam was removed and the proton flux on the \(K_S\) target was greatly increased. A 24 mm platinum absorber was placed after the Be target to reduce the photon flux in the neutral beam. Using this beam line, together with the rare \(K_S\), NA48/1 produced many Hyperons decays results from which the first measurement of \(\Xi^0 \to \Lambda e^+e^-\) will be discussed. Since 2003 the neutral beams were replaced by simultaneous \(K^+\) and \(K^-\) beams for the NA48/2 experiment. The momentum \((60 \pm 3)\) GeV/c was formed symmetrically for \(K^+\) and \(K^-\) in the first achromat (see Fig. 1), in which the two beams were split in the vertical plane. In the second achromat were placed two of the three stations of the Kaon beam spectrometer (KABES). The beams followed the same path in the decay volume, comprised in a 114 m long cylindrical vacuum tank. The beam axes coincided to 1 mm, while their lateral size is about 1 cm. The results on charged 3\(\pi\) asymmetry, \(K^\pm \to \pi^\pm\pi^0\gamma\), \(K^\pm \to \pi^\pm e^+e^-\gamma\), semileptonic decays, and \(K^\pm \to e^\pm\nu/K^\mp \to \mu^\pm\nu\) exploit this last beam configuration.

3 Charge asymmetry in \(K^\pm \to \pi^\pm\pi^\mp\pi^\pm\) and \(K^\pm \to \pi^\pm\pi^0\pi^0\)

The standard phenomenological description of \(K_{3\pi}\) decays is made in terms of the bi-dimensional Dalitz plot variables \(u\) and \(v\), related respectively to the energy sharing to the "odd" pion (charge opposite with respect to the other two) and among the two "even" pions:

\[
u = \frac{s_2 - s_1}{m^2_{\pi^0}}
\]
Figure 2: Experimental status of asymmetry in $K^\pm \rightarrow \pi^\pm \pi^+ \pi^-$.

Figure 3: Experimental status of asymmetry in $K^\pm \rightarrow \pi^\pm \pi^0 \pi^0$.

where $s_i = (P_{K_i} - P_{\pi_i})^2$, $i = 1, 2, 3$, $s_0 = (s_1 + s_2 + s_3)/3$ being $P_{K_i}$ and $P_{\pi_i}$ the Kaon and Pion four-momenta, the indexes $i = 1, 2$ correspond to the two identical pions and the index $i = 3$ to the pion of opposite charge. Using those two variables the matrix element can be parametrized as a polynomial expansion with slopes to be measured in experiments:

$$|M(u, v)|^2 = 1 + gu + hu^2 + ku^2$$

(2)

where $g(\pi^+\pi^+\pi^-) = 0.2154 \pm 0.0035$, $g(\pi^+\pi^0\pi^0) = 0.638 \pm 0.020$ and $|h|, |k| \ll |g|$. Any difference in the slopes parameters $g^+$ and $g^-$ for positive and negative kaon, is a clear manifestation of direct CP violation. Usually the asymmetry parameter is defined by the corresponding slope asymmetry:

$$A_g = \frac{g^+ - g^-}{g^+ + g^-} \approx \frac{\Delta g}{2g}$$

(3)

SM predictions for that asymmetry vary between few $10^{-6}$ to few $10^{-5}$ while in models beyond the SM it can be enhanced up to the low $10^{-4}$ region. The parameter $\Delta g$ can be extracted experimentally by comparing the reconstructed $u$ spectra of $K^+$ and $K^-$ decays ($N^+(u)$ and $N^-(u)$): in the $K^\pm \rightarrow \pi^\pm \pi^\pm \pi^\pm$ case the ratio $R(u) = N^+(u)/N^-(u)$ is in good approximation proportional to $(1 + \Delta g \cdot u)$, so $\Delta g$ can be extracted from a linear fit. In the past years, several experiments have searched for the CP violating slope asymmetry in both $K^\pm \rightarrow \pi^\pm \pi^\pm \pi^\pm$ and $K^\pm \rightarrow \pi^\pm \pi^0 \pi^0$ decay modes by collecting samples of $K^+$ and $K^-$ decays. These measurements set upper limits on $A_g$ at the level of a few $10^{-3}$, limited by systematic uncertainties.

In NA48/2 using simultaneous and collinear $K^+$ and $K^-$ beams and reversing all the magnets polarities during the data taking allows a charge symmetrization of the experimental conditions lowering the systematic uncertainties to few in $10^{-4}$. Data collected over a period with all the four possible magnets setup configurations represent a "SuperSample", which is treated as an independent and self-consistent set of data for asymmetry measurement. To measure the charge asymmetry the following "quadruple ratio" composed as a product of four $R(u) = N^+(u)/N^-(u)$ ratios with opposite kaon sign, is used: $R4(u) = R_{US}(u) \cdot R_{DJ}(u) \cdot R_{DS}(u) \cdot R_{DJS}(u)$ where the indices $U$ (D) denote the beam line polarities corresponding to $K^+$ passing along the Upper (Lower) path in the achromats, respectively, while the indices $S$ (J) represent spectrometer magnet polarities corresponding to the "even" pions being deflected to negative (positive) $x$, i.e. towards the Salve (Jura) mountains, respectively.
3.1 $K^\pm \to \pi^\pm \pi^\mp \pi^\pm$ asymmetry final result

A very simple selection based on the measured momenta of the three charged track requires the compatibility of vertexes computed tracking back the direction of couples of tracks, reconstructed Kaon momentum compatible with the beam momentum, and the kaon reconstructed mass being in the range $M_K (PDG) \pm 9$ MeV. Using the full 2003-2004 data sample the selection leaves a practically background free sample of $3.82 \cdot 10^9$ as $K^\pm \to \pi^\pm \pi^\mp \pi^\pm$ is the dominant three-track decay mode.

Many sources of systematic effects were investigated such as the fine alignment of the spectrometer, the geometrical acceptance seen by the two beams, the dependence on the way the $u$ variable is calculated or the fitting limits, effects due to uncertainty on the knowledge of the magnetic fields, pile-up effects, inhomogenities in the spectrometer alignment and trigger efficiencies. After applying the above corrections, $A_g$ is extracted by fitting the quadruple ratio of the $u$ spectra for each SuperSample.

The difference in $K^\pm \to \pi^\pm \pi^0 \pi^\pm$ Dalitz plot slope parameter is found to be:

$$\Delta g^C = (0.7 \pm 0.7_{stat} \pm 0.4_{trig} \pm 0.5_{syst}) \cdot 10^{-4}$$

(4)

Converted to the direct CP violating charge asymmetry using the value of the Dalitz plot slope $g^C = -0.21134 \pm 0.00017$ recently measured by the NA48/2:

$$A_g^C = (-1.5 \pm 1.5_{stat} \pm 0.9_{trig} \pm 1.1_{syst}) \cdot 10^{-4} = (-1.5 \pm 2.1) \cdot 10^{-4}$$

(5)

which shows no CP violation in agreement with the SM prediction. The result has $\sim 20$ times better precision than the best measurement before NA48/2, see Fig. 2 and the precision is still limited mainly by the available statistics.

3.2 $K^\pm \to \pi^\pm \pi^0 \pi^0$ asymmetry final result

The event selection requires 1 charged track and at least 4 in time photon clusters fulfilling geometrical and quality requirements. For each selected event, the decay is reconstructed as follows. Assuming that each pair i, k of LKr clusters (i, k =1,2,3,4) originates from a $\pi^0$ decay, the distance $D_{ik}$ between the $\pi^0$ decay vertex position along the z axis and the front plane of the LKr is calculated. Among all photon pairs, the two with the smallest $D_{ik}$ difference are selected as the best combination consistent with the hypothesis of two $\pi^0$ mesons originating from $K^\pm \to \pi^\pm \pi^0 \pi^0$ decay. Using the mean of the two vertex positions the kaon mass and momenta are reconstructed. Further event selection requires the $\pi^\pm \pi^0 \pi^0$ invariant mass to differ from the nominal $K^{pm}$ mass by less than 6 MeV/c$^2$, and the reconstructed kaon momentum to be between 54 and 66 GeV/c.

The above requirements applied to the full 2003 and 2004 data sample lead to the final sample of $9.13 \cdot 10^7$ events. The background is negligible for the applied mass cut. It can be seen that the kinematic variable $u$ can be computed using only the $\pi^0 \pi^0$ invariant mass. Thus a measurement of $u$ uses the information from the LKr only, not involving the DCH data. This provides a certain charge symmetry of the procedure, as the LKr is a charge blind detector, except for effects of small differences between $\pi^+$ and $\pi^-$ interaction characteristics.

The difference in the linear slope parameters of $K^+$ and $K^-$ decays into $\pi^\pm \pi 0\pi 0$ was measured to be

$$\Delta g^n = (2.2 \pm 2.1_{stat} \pm 0.6_{syst}) \cdot 10^{-4}$$

(6)

The corresponding direct CP violating asymmetry obtained using the nominal value of the linear slope parameter $g^n = 0.626 \pm 0.007^2$ is

$$A_g^n = (1.8 \pm 1.7_{stat} \pm 0.5_{syst}) \cdot 10^{-4} = (1.8 \pm 1.8) \cdot 10^{-4}$$

(7)
As for $K^\pm \rightarrow \pi^\pm \pi^0 \pi^0$ an order of magnitude improvements has been achieved with respect to previous measurements, see Fig. 3, and no CP violation has been observed according to current standard model predictions.

4 The radiative decay $K^\pm \rightarrow \pi^\pm \pi^0 \gamma$

The decay channel $K^\pm \rightarrow \pi^\pm \pi^0 \gamma$ is one of the most interesting and important channels for studying the low energy structure of the QCD, in fact it has been shown to be one of the most sensitive check for the Chiral Anomaly. The total amplitude of $K^\pm \rightarrow \pi^\pm \pi^0 \gamma$ decay is the sum of two terms: the inner bremsstrahlung (IB) associated with the decay $K^\pm \rightarrow \pi^\pm \pi^0$, in which the photon is emitted from the outgoing charged pion, and the direct emission (DE) in which the photon is emitted from one of the intermediate states of the decay itself. Although, due to the dominant IB, the DE component is very difficult to observe it can be isolated kinematically. As the $K^\pm \rightarrow \pi^\pm \pi^0$ decay is suppressed by the $\Delta I = 1/2$ rule, also the IB will similarly suppressed. This feature could then enhance the DE. The DE term can occur through both electric (E) and magnetic (M) dipole transitions. While the magnetic part, can be evaluated using the Wess-Zumino-Witten functional, describing the reducible anomaly, there is no definite prediction from ChPT on the electric transition, whose amplitude depends on undetermined constants. The electric contribution is extremely interesting since it interferes (INT) with the IB amplitude therefore it may be distinguished from the magnetic, which does not.

In the $K^\pm \rightarrow \pi^\pm \pi^0 \gamma$ decay IB, INT, and DE components can be separated kinematically using the Lorentz invariant variable $W$ which is defined as follows:

$$ W^2 = \frac{(P_K^* \cdot P_x^*)(P_{\pi}^* \cdot P_{\gamma}^*)}{(m_K m_{\pi})^2} \quad (8) $$

with $P_x^*$ the 4-momentum of the $x$ particle and $\gamma$ the radiative one. The decay rate depends only on $T_{\pi}^*$, energy of the pion in the Kaon rest frame, and $W$. Integrating over $T_{\pi}^*$ an expression
that splits the different contributions into terms with different powers of $W$ can be obtained:

$$\frac{dI^\pm}{dW} \approx \left( \frac{dI^\pm}{dW} \right)_{IB} \left[ 1 + 2 \left( \frac{m_{\pi}}{m_K} \right)^2 W^2 |E| \cos((\delta_1 - \delta_0) \pm \phi) + \left( \frac{m_{\pi}}{m_K} \right)^4 W^4 (|E|^2 + |M|^2) \right]$$ (9)

where $|E|$ and $|M|$ represent electric and magnetic transitions, while $\phi$ is an unknown phase responsible for CP violation. The three terms represent IB, INT and DE contribution respectively. Recent paper by D’Ambrosio and Cappiello suggest the presence of a form factor in the DE term, not yet included in the present analysis, that complicates a little bit the very simple parametrization in Eq. 9.

The IB component has been measured since the seventies by Abrams et al.\textsuperscript{4} achieving a good agreement with solid QED theoretical predictions. The experimental measurement of the fractions of DE and INT is affected by very dangerous BG sources, such as $K^\pm \to \pi^\pm \pi^0$ and $K^\pm \to \pi^\pm \pi^0 \pi^0$ decays, suppressed in a kinematically BG free region, $T_x^*$ in the range 55-90 MeV. Moreover a good measurement requires a very good reconstruction of both charged and neutral particles 4-momenta. The present experimental knowledge about the decay is limited to the DE component while the INT has been never observed. The most recent results are summarized in Tab. 1. All of them have been obtained in the $T_x^*$ region 55-90 MeV and setting INT=0 in the fitting procedure.

<table>
<thead>
<tr>
<th>exp.</th>
<th>year</th>
<th>#events</th>
<th>$BR(DE) \cdot 10^{-6}$</th>
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<tr>
<td>E787\textsuperscript{5}</td>
<td>2000</td>
<td>20K</td>
<td>4.7 ± 0.8 ± 0.3</td>
</tr>
<tr>
<td>E470\textsuperscript{6}</td>
<td>2003</td>
<td>4.5K</td>
<td>3.2 ± 1.3 ± 1.0</td>
</tr>
<tr>
<td>E787\textsuperscript{7}</td>
<td>2005</td>
<td>20K</td>
<td>3.5 ± 0.6 ± 0.35</td>
</tr>
<tr>
<td>E470\textsuperscript{8}</td>
<td>2005</td>
<td>10K</td>
<td>3.8 ± 0.8 ± 0.7</td>
</tr>
<tr>
<td>ISTRA+\textsuperscript{9}</td>
<td>2006</td>
<td>930</td>
<td>3.7 ± 3.9 ± 1.0</td>
</tr>
</tbody>
</table>

Table 1: The $K^\pm \to \pi^\pm \pi^0 \gamma$ experimental results

The errors on the measurements are quite high and the agreement in not so good. Moreover the assumption of vanishing interference term, known to exist by theory, is based on the following measurements by both E787\textsuperscript{5} and E470\textsuperscript{6}:

\begin{align*}
INT &= (-0.4 \pm 1.6)\% \quad E787 \quad (10) \\
INT &= (-0.58_{-0.83}^{+0.01})\% \quad E470 \quad (11)
\end{align*}

providing very weak constraints on the true value of the INT term.

The main BG sources are $K^\pm \to \pi^\pm \pi^0$ and $K^\pm \to \pi^\pm \pi^0 \pi^0$. The first decay needs an accidental photon or an hadronic extra cluster to mimic the signal final state, while the second a lost or 2 fused gamma. The selection aims to suppress the contribution of them both to less than 1% of the DE component. The rejection of $K^\pm \to \pi^\pm \pi^0$ relaxes on the $T_x^*$ cut. The request $T_x^*$ lower than 80 MeV allows to reject $K^\pm \to \pi^\pm \pi^0$ and a part of the IB spectrum of $K^\pm \to \pi^\pm \pi^0 \gamma$ only, including region 0-55 MeV very rich of DE and INT events. The upper cut at 80 MeV is due to trigger reasons. To suppress $K^\pm \to \pi^\pm \pi^0 \pi^0$ BG the very good kaon mass resolution, (2.2 MeV), and the identification of fused gamma events constraints have been used. In Fig. 4 the data kaon mass spectrum is compared with the sum of $K^\pm \to \pi^\pm \pi^0 \gamma$ and $K^\pm \to \pi^\pm \pi^0 \pi^0$ MC. The figure shows that the background contribution is very low and that can be explained in term of $K^\pm \to \pi^\pm \pi^0 \pi^0$ only.

A very important issue in the measurement of DE and INT is to identify the radiative $\gamma$ among the 3 available. A dedicated set of cuts, based on the agreement of the vertex evaluated using the pion and kaon tracks, and the one evaluated pairing the $\gamma$'s to form a $\pi^0$, has been
implemented. Using this cuts a misidentification probabilities, computed using MC simulation, of the order of the permille for all the components has been achieved. A total of \( \sim 220000 \) candidate events survived all selection cuts in the region of \( T^*_\pi \) 0-80 MeV using only the first 3 super sample of the 2003 data set.

The extraction of the fractions of DE and INT relies on the fact that different components show quite different \( W \) distributions. An extended maximum likelihood technique, assigning weights to MC \( W \) distributions of the 3 components to reproduce data spectrum, has been employed to get the fractions. Many checks to verify the result were made concerning the \( \gamma \) energies reconstructed in the calorimeter, trigger efficiencies and BG contribution. The systematic uncertainties are dominated by the trigger.

To extract the fractions the fit has been performed in the interval 0.2-0.9 in the \( W \) variable and with a minimum gamma energy of 5 GeV using only 124K events from the total sample. After correcting for different acceptances we get, in the region \( 0 < T^*_\pi < 80 \text{MeV} \), the following preliminary values for the fractions of DE and INT with respect to IB:

\[
\text{Frac}(\text{DE}) = (3.35 \pm 0.35_{\text{stat}} \pm 0.25_{\text{syst}})\% \\
\text{Frac}(\text{INT}) = (-2.67 \pm 0.81_{\text{stat}} \pm 0.73_{\text{syst}})\%
\]

This is the first measurement of a non vanishing interference term in the \( K^\pm \rightarrow \pi^\pm \pi^0 \gamma \gamma \) decay. The contour plot in Fig. 5 shows the very high correlation of the two contributions \( \rho = -0.92 \). Major improvements in the size of the data sample are foreseen using the full 2003-2004 data set.

5 First observation of the decay \( K^\pm \rightarrow \pi^\pm e^+ e^- \gamma \)

The decay \( K^\pm \rightarrow \pi^\pm e^+ e^- \gamma \) has a kinematic very similar to the corresponding \( K \rightarrow \pi^\pm \gamma \gamma \). One of the photons internally converts into a pair of electrons. The branching ratio can be naively estimated using the following relation:

\[
\text{BR}(K^\pm \rightarrow \pi^\pm e^+ e^- \gamma) = K \rightarrow \pi^\pm \gamma \gamma \cdot 2\alpha = 1.6 \cdot 10^{-8}
\]

Both decays are described by chiral perturbation theory. Their lowest order terms are \( O(p^4) \), leading to a characteristic signature in the \( e^+ e^- \) mass of the decay. Model dependent theoretical
estimates based on chiral perturbation theory predicts the BR to be in the range $^{11}\,$:

\[ BR(K^+ \to \pi^+e^+\gamma) = (0.9 - 1.6) \cdot 10^{-8} \]  \hspace{1cm} (15)

In the paper the differential distribution on $m_{\text{exc}}$, very similar to the $K^\pm \to \pi^\pm\gamma\gamma$ one, is also given in term of the same $\hat{\chi}_c$ parameter.

NA48/2 experiment observed for the first time the radiative decay $K^\pm \to \pi^\pm e^+\gamma$. 92 candidates were selected, with $1 \pm 1$ accidental background and $5.1 \pm 1.7$ misidentification background. The main source of BG is the $K \to \pi^+\pi^0\gamma$ with a lost $\gamma$. In Fig. 6 the reconstructed Kaon invariant mass is shown. The crosses represent data while the filled distributions represent different simulated background contributions.

By using $K^\pm \to \pi^\pm\pi^0$ as normalization channel the branching ratio was preliminary estimated to be

\[ BR(K^\pm \to \pi^\pm e^+\gamma) = (1.27 \pm 0.14_{\text{stat}} \pm 0.05_{\text{syst}}) \cdot 10^{-8} \]  \hspace{1cm} (16)

Unfortunately due to lack of statistics the value of $\hat{\chi}_c$ is not accessible to this analysis.

6 Semileptonic Decays results

The branching ratios of semileptonic kaon decays are needed to determine $|V_{US}|$ element in the CKM matrix. In addition $\Gamma(K e3)/\Gamma(K \mu3)$ is a function of the slope parameters of the form factors, which can be used for consistency check under the assumption of lepton universality.

A special run dedicated to collect semileptonic decays has been performed during 2003 data taking of NA48/2. Approximately 56000 $K_3^\leftrightarrow$, 31000 $K_e3$, 49000 $K_\mu3$, 28000 $K_\mu3$, 462000 $K_{2\pi}$, and 256000 $K_{2\pi}$ decays were selected for the measurement. The ratios of decay widths, combining $K^+$ and $K^-$, are:

\[ \Gamma(K e3)/\Gamma(K 2\pi) = 0.2470 \pm 0.0009_{\text{stat}} \pm 0.0004_{\text{syst}} \]  \hspace{1cm} (17)

\[ \Gamma(K_\mu3)/\Gamma(K 2\pi) = 0.1637 \pm 0.0006_{\text{stat}} \pm 0.0003_{\text{syst}} \]  \hspace{1cm} (18)

\[ \Gamma(K_\mu3)/\Gamma(K e3) = 0.663 \pm 0.003_{\text{stat}} \pm 0.001_{\text{syst}} \]  \hspace{1cm} (19)

Using the PDG '06$^2$ value for the $K_{2\pi}$ branching fraction, 0.2092 ± 0.0012, those for the semileptonic decays have been computed to be:

\[ BR(K e3) = (5.167 \pm 0.019_{\text{stat}} \pm 0.008_{\text{syst}} \pm 0.030_{\text{norm}})\% \]  \hspace{1cm} (20)

\[ BR(K_\mu3) = (3.425 \pm 0.013_{\text{stat}} \pm 0.006_{\text{syst}} \pm 0.020_{\text{norm}})\% \]  \hspace{1cm} (21)
The uncertainty is dominated by the existing data for the BR($K_{2\pi}$) used as normalization. By using the measured values for the vector and the scalar form factors, and assuming $e-\mu$ universality, the ratio $\Gamma(Ke3)/\Gamma(K\mu3)$ can be expressed as:

\[
\Gamma(Ke3)/\Gamma(K\mu3) = \frac{0.645 + 2.087\lambda_+ + 1.464\lambda_0 + 3.375\lambda_1^2 + 2.573\lambda_3^2}{1 + 3.457\lambda_+ + 4.783\lambda_1^2}
\]

This relationship can be used to crosscheck the form factors measurements. The value measured by NA48/2 for the ratio $0.663 \pm 0.003_{\text{stat}} \pm 0.001_{\text{sys}}$ is consistent with KEK E246 and PDG '06. Given the values of the semileptonic decays BRs the product $|V_{US}|f_+(0)$ can be extracted using the following formula:

\[
\frac{BR(K_{3\pi})}{\tau_K} = \frac{C_K^2 G_F^2 m_K^5}{192\pi^3 S_{EW}} |V_{US}|^2 |f_+(0)|^2 I_K^l(\lambda_0)(1 + \delta_{SU(2)}^l + \delta_{EM}^l)
\]

we get:

\[
\text{From } Ke3 : \quad |V_{US}|f_+(0) = 0.2193 \pm 0.0012
\]

\[
\text{From } K\mu3 : \quad |V_{US}|f_+(0) = 0.2177 \pm 0.0013
\]

in which the errors are dominated by the uncertainties of the external inputs needed for the calculation in Eq. 23. Combining the results from both modes, assuming lepton universality and taking the value of $f_+(0)$ for neutral kaons, the obtained $|V_{US}|$ element is:

\[
|V_{US}| = 0.2185 \pm 0.0023
\]

which is consistent with CKM matrix unitarity predictions. For detailed description of the analysis see\textsuperscript{12}.

7 The ratio $K^\pm \rightarrow e^\pm \nu/K^\pm \rightarrow \mu^\pm \nu$

The measurement of the ratio $R_K = K^\pm \rightarrow e^\pm \nu/K^\pm \rightarrow \mu^\pm \nu$ between the decay rates of $K^\pm \rightarrow e^\pm \nu$ and $K^\pm \rightarrow \mu^\pm \nu$ is a sensitive test of lepton universality and of V-A structure of the weak interactions. In fact, while in the BRs of $K \rightarrow \nu$ the theoretical uncertainties are at the percent level, due to $f_K$, in the ratio of the electronic and muonic decay modes, the hadronic uncertainties cancel to a very large extent. As a result, the SM predictions for $R_K$ is known with excellent accuracy and this makes it possible to fully exploit the great experimental resolutions in the ratio to constrain new physics effects. In the standard model $R_K$ is given by:

\[
R_K = \frac{\Gamma(K \rightarrow e\nu(\gamma))}{\Gamma(K \rightarrow \mu\nu(\gamma))} = \frac{m_e^2}{m_\mu^2} \left( \frac{m_K^2 - m_e^2}{m_K^2 - m_\mu^2} \right)^2 (1 + \delta R_K) = (0.2472 \pm 0.001) \cdot 10^{-5}
\]
where $\delta R_K$ is due to the difference in the radiative corrections. The experimental value given by PDG$^2$ is $R_K = (0.245 \pm 0.11) \cdot 10^{-5}$ less precise by far.

Recent paper by Masiero, Paradisi, and Petronzio$^3$ studies possible contribution of Lepton Flavor Violation (LFV) mechanisms introduced by SUSY models to the value of $R_K$. The one involved in the calculation arise from a charged Higgs exchange through the diagram in Fig. 10.

The result of the calculation give the above enhancement for the $R_K$:

$$R_{K^{\text{LFV}}}^{\text{SM}} = R_{K^{\text{SM}}}^{\text{SM}} \cdot \left[1 + \left(\frac{m_K^2}{M_H^2}\right) \left(\frac{m^2}{m_\mu^2}\right) |\Delta R|^2 \tan^6 \beta \right]$$

(28)

In the large $\tan \beta$ regime ($\sim 40$) and with a relatively heavy $H^\pm$ ($M_H = 500$ GeV), it is possible to reach deviations from SM $R_K$ at the percent level, thanks to LFV enhancements arising in the SUSY model.

Na48/2 has collected the world largest sample of $K^{\pm} \rightarrow e^{\pm}\nu$ and is therefore able to perform a very precise measurement of $R_K$. The measurement of the ratio allows the cancelation of acceptance and trigger effects common to both type of events. The selection is very similar except for the particle identification part, based on the reconstruction of the ratio E/P. In the 2003 run 5239 $K^{\pm} \rightarrow e^{\pm}\nu$ were selected, by a downscaled trigger, with $> 14\%$ background mainly from $K^{\pm} \rightarrow \mu^{\pm}\nu$. In fact in cases in which the $\mu$ produces a very high energy bremsstrahlung photon, the value of E/P is very near to one, and the event is misidentified as an electron type events. Even if the BG from this source is subtracted in the ratio calculation it is still the major source of systematic uncertainties. The preliminary result obtained with the 2003 data set only is:

$$R_K = (2.416 \pm 0.043_{\text{stat}} \pm 0.024_{\text{syst}}) \cdot 10^{-5}$$

(29)

The estimations yield that the combined 2003 and 2004 result will not be enough to obtain a total error smaller than 1%. A dedicated 2007 run will be performed during upcoming summer. The conservative estimation for the error, which will be reached in $R_K$ measurement is 0.7$\%$. In the future experiment P326 a per mill uncertainty could be reached, due to better particle identification provided by the RICH detector.

8 First observation of $\Xi^0 \rightarrow \Lambda e^+e^-$

In the 2002 NA48/1 run the weak radiative decay $\Xi^0 \rightarrow \Lambda e^+e^-$ was detected for the first time. 412 candidates were selected with 15 background events Fig. 11. The obtained branching fraction is:

$$BR(\Xi^0 \rightarrow \Lambda e^+e^-) = (7.7 \pm 0.5_{\text{stat}} \pm 0.4_{\text{syst}}) \cdot 10^{-6}$$

(30)
is consistent with inner bremsstrahlung-like $e^+e^-$ production mechanism. The decay parameter $\alpha_{\Xi\Lambda\Lambda\Lambda}$ can be measured from the angular distribution:

$$\frac{dN}{d\cos \theta_{p\Xi}} = \frac{N}{2} \left( 1 - \alpha_{\Xi\Lambda\Lambda\Lambda} \cos \theta_{p\Xi} \right)$$  \hspace{1cm} (31)$$

where $\theta_{p\Xi}$ is the angle between the proton from $\Lambda \to p\pi$ decay relative to the $\Xi^0$ line of flight in the $\Lambda$ rest frame and where $\alpha_\Lambda$ is the asymmetry parameter for the decay $\Lambda \to p\pi^-$. The obtained value $\alpha_{\Xi\Lambda\Lambda\Lambda} = -0.8 \pm 0.2$ is consistent with the latest published value of the decay asymmetry parameter for $\Xi \to \Lambda\gamma$. Detail on the analysis can be found in 14.

9 $K_L \to \pi^\pm \mu^\mp \nu$ form factors

$K\bar{\Lambda}$ decays provide the cleanest way to extract $|V_{US}|$ element in the CKM matrix. Recent calculations in the framework of $\chi PT$ show how the vector form factor at zero momentum transfer, $f_+(0)$, can be constrained experimentally from the slope and curvature of the scalar form factor $f_0$ of the $K_L \to \pi^\pm \mu^\mp \nu$ decay. In addition, these form factors are needed to calculate the phase space integrals, which are used in $|V_{US}|$ determination. Approximately $2.6 \cdot 10^6 K_L\bar{\Lambda}$ decays were selected from a dedicated 1999 minimum bias run of NA48. By studying the Dalitz plot density, the following slopes for the vector and the scalar form factors were obtained:

$$\lambda'_+ = (20.5 \pm 2.2_{\text{stat}} \pm 2.4_{\text{syst}}) \cdot 10^{-3}$$  \hspace{1cm} (32)$$

$$\lambda'_0 = (2.6 \pm 0.9_{\text{stat}} \pm 1.0_{\text{syst}}) \cdot 10^{-3}$$  \hspace{1cm} (33)$$

$$\lambda_0 = (9.5 \pm 1.1_{\text{stat}} \pm 0.8_{\text{syst}}) \cdot 10^{-3}$$  \hspace{1cm} (34)$$

The results show the presence of a quadratic term in the expansion of the vector form factor in agreement with other recent measurements. A comparison between the results of the quadratic fits as reported by the recent experiments is presented in Fig. 12. The results obtained with linear fit are

$$\lambda'_+ = (26.7 \pm 0.6_{\text{stat}} \pm 0.8_{\text{syst}}) \cdot 10^{-3}$$  \hspace{1cm} (35)$$

$$\lambda_0 = (11.7 \pm 0.7_{\text{stat}} \pm 1.0_{\text{syst}}) \cdot 10^{-3}$$  \hspace{1cm} (36)$$
The value for $\lambda_+$ is well compatible with the recent KTeV measurement, while $\lambda_0$ is shifted towards lower values. Details on NA48 $K\mu3$ measurement can be found in\cite{15}.

References

7. T. Tsunemi, “New Results on $K^+ \rightarrow \pi^+\pi^0\gamma$ from E787”, talk given at Kaon 2005 Chicago, June 2005.
RECENT RESULTS ON KAON DECAYS FROM KLOE AT DAΦNE

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The analysis of the full data sample of about 2.5 fb$^{-1}$ collected by the KLOE experiment at DADFNE is in progress. New results on $K_S^+ \rightarrow e^+e^-$, $K_S^+ \rightarrow \gamma\gamma$, $K_L \rightarrow \pi\nu\bar{\nu}$, and $K_{\mu3}$ form factor are presented. The most recent KLOE results on $V_{us}$, CPT invariance and QM tests are reviewed.

1 The KLOE experiment at DADFNE

The Frascati $\phi$-factory, DAΦNE, is an $e^+e^-$ collider working at a center of mass energy of $\sqrt{s} \approx 1020$ MeV, corresponding to the peak of the $\phi$ resonance. The $\phi$ production cross section is $\sim 3\mu$b, and the $\phi \rightarrow K^0\bar{K}^0$ decay has a branching fraction of 34%. The beams collide at the interaction point (IP) with a crossing angle $\theta_x \approx 25$ mrad, therefore $\phi$s are produced with a small momentum of $\sim 12.5$ MeV in the horizontal plane. The typical sizes of the beam are:

$\sigma_x = 0.2\text{ cm}; \sigma_y = 20 \mu\text{m}; \sigma_z = 3 \text{ cm}$.

The KLOE detector consists mainly of a large volume drift chamber surrounded by an electromagnetic calorimeter. A superconducting coil provides a 0.52 T solenoidal magnetic field.

The fine sampling lead-scintillating fiber calorimeter (ECM) consists of a barrel and two end-caps, and has solid angle coverage of 98%. Photon energies and arrival times are measured with resolutions $\sigma_E/E = 5.7\%/\sqrt{E(\text{GeV})}$ and $\sigma_t = 54\text{ ps}/\sqrt{E(\text{GeV})} \oplus 50\text{ ps}$, respectively. Photon entry points are determined with an accuracy $\sigma_z \sim 1 \text{ cm}/\sqrt{E(\text{GeV})}$ along the fibers and $\sigma_\perp \sim 1 \text{ cm}$ in the transverse direction.

The tracking detector is a 4 m diameter and 3.3 m long cylindrical drift chamber (DCH) with a total of $\sim 52000$ wires, of which $\sim 12000$ are sense wires. In order to minimize multiple scattering and $K_L$ regeneration and to maximize detection efficiency of low energy photons, the chamber works with a helium based gas mixture and its walls are made of light materials (mostly...
carbon fiber composites). The momentum resolution for tracks produced at large polar angle is \( \sigma_p/p \leq 0.4\% \). Vertices are reconstructed with a resolution of \( \sim 3 \) mm.

At a \( \phi \)-factory the \( \phi \rightarrow K^0\bar{K}^0 \) decay produces the neutral kaon pair in a coherent quantum state with \( J^{PC} = 1^- \):

\[
|i\rangle = \frac{1}{\sqrt{2}} \{ |K^0\rangle |\bar{K}^0\rangle - |\bar{K}^0\rangle |K^0\rangle \} = \frac{N}{\sqrt{2}} \{ |K_S\rangle |K_L\rangle - |K_L\rangle |K_S\rangle \}
\]

where \( N = (1 + |\epsilon|^2)/(1 - \epsilon^2) \simeq 1 \) is a normalization factor.

The detection of a kaon at large (small) times tags a \( K_S \) (\( K_L \)) in the opposite direction. At KLOE a \( K_S \) is tagged by identifying the interaction of the \( K_L \) in the calorimeter (\( K_L \)-crash). In fact about 50% of the produced \( K_L \)'s in \( \phi \rightarrow K_SK_L \) events reach the calorimeter before decaying; their associated interactions are identified by a high energy, neutral and delayed deposit in the calorimeter, i.e. not associated to any charged track in the event, and delayed of \( \sim 30 \) ns (as at KLOE \( K_L \)'s have a velocity \( \beta \sim 0.22 \)) with respect to a photon coming from the interaction region. Pure \( K_S \) samples have been selected exploiting this tagging technique.

A \( K_L \) is tagged by detecting a \( K_S \rightarrow \pi^+\pi^- \) decay near the IP; the invariant mass from the momenta of the two pion tracks is reconstructed with a resolution of \( \sim 1 \) GeV, thus allowing the selection of a clean \( K_L \) sample.

KLOE completed the data taking in March 2006 with a total integrated luminosity of \( \sim 2.5 \) fb\(^{-1}\), corresponding to \( \sim 7.5 \times 10^{-9} \) \( \phi \)-mesons produced.

2 New KLOE results on neutral kaon decays

2.1 \( K_S \rightarrow e^+e^- \)

The SM prediction of the branching ratio of \( K_S \rightarrow e^+e^- \) decay is rather low (\( \text{BR} = 1.6 \times 10^{-15} \)), but quite precise\(^3\), leaving room for possible new physics effects to be detected. A data sample corresponding to 1.3 fb\(^{-1}\) has been analyzed; events are selected requiring the presence of a \( K_S \) decay, tagged by the detection of a \( K_L \)-crash, and two charged tracks coming from the IP, with an invariant mass (in the \( e^+e^- \) hypothesis) \( M_{\text{inv}} \) greater than 420 MeV. A \( \chi^2 \)-like variable is built, based on the measured time of flights of the two particles, \( E/p \), and the transverse distance between the track impact point on the calorimeter and the closest calorimeter cluster.

The search for the signal is performed inside a signal box in the \( \chi^2-M_{\text{inv}} \) plane, whose definition is optimized with a Monte Carlo (MC) simulation study: \( 492 \leq M_{\text{inv}} \leq 504 \) MeV and \( \chi^2 \leq 20 \). These cuts reject almost all the events due to the background processes \( K_S \rightarrow \pi^+\pi^- \rightarrow \mu\pi, K_S \rightarrow \pi^+\pi^- \), and \( \phi \rightarrow \pi^+\pi^- \), while retaining 55.8% of the signal. The selection is inclusive for radiated photons with energy in the kaon rest frame \( E_\gamma \leq 6 \) MeV. We observe \( N = 3 \) events inside the box, with an expected background of \( 7.1 \pm 3.6 \) events, corresponding to an upper limit of 4.3 events at 90% c.l.; after normalization to \( K_S \rightarrow \pi^+\pi^- \) events, we obtain a preliminary result for the branching fraction:

\[
\text{BR}(K_S \rightarrow e^+e^-(\gamma); E_\gamma \leq 6 \text{ MeV}) < 2.1 \times 10^{-8}
\]

at 90% c.l., improving the previous limit\(^4\) by almost an order of magnitude.

2.2 \( K_S \rightarrow \gamma\gamma \)

The measurement of \( \text{BR}(K_S \rightarrow \gamma\gamma) \) is an important test of chiral perturbation theory, as discussed in Ref.\(^5\). A data sample corresponding to 1.6 fb\(^{-1}\) has been analyzed; events are selected requiring the presence of a \( K_S \) decay, tagged by the detection of a \( K_L \)-crash, and two and only two photons with an energy greater than 7 MeV, back-to-back in the center of mass
system of $K_S (\cos(\theta^*_\gamma) \leq -0.95)$, and coming from the IP ($T_\gamma - R/e \simeq 0$). A kinematic fit constrains the time, momentum, and invariant mass $M_{\gamma\gamma}$ of the two photons. In order to further reduce the copious background due to $K_S \rightarrow 2\pi^0$ with two undetected photons, a veto on the signal of two small calorimeters surrounding the focusing quadrupoles near the IP is applied. The background due to $K_L \rightarrow \gamma\gamma$ decays is absent due to the high purity of the $K_S$ sample. The overall efficiency on the signal is $\sim 52\%$. Finally, the signal counts are obtained by fitting the $M_{\gamma\gamma}$ and $\cos(\theta^*_\gamma)$ bi-dimensional distribution with signal and background distributions obtained from MC. The reconstructed energy scale is well kept under control by comparing data samples of $K_S \rightarrow 2\pi^0$ and $K_L \rightarrow \gamma\gamma$ events with MC. The preliminary result is

$$\text{BR}(K_S \rightarrow \gamma\gamma) = (2.35 \pm 0.14) \times 10^{-6}$$

in agreement with $O(p^4)$ chiral perturbation theory calculation, and not confirming the discrepancy of $\sim 30\%$ found by the NA48 collaboration\(^6\).

2.3 $K_L \rightarrow \pi e\nu\gamma$

Radiative effects play an important role in kaon semileptonic decays. Both inner bremsstrahlung (IB) and structure dependent (SD) amplitudes contribute to the $K_L \rightarrow \pi e\nu\gamma$ process, as discussed in Ref.\(^7\). A data sample corresponding to $\sim 330$ pb$^{-1}$ has been analyzed; inclusive selection of $K_L \rightarrow \pi e\nu(\gamma)$ events requires a $K_L$ of known momentum and direction, tagged by $K_S \rightarrow \pi^+\pi^-$ decay near the IP. In a fiducial volume extending for $\sim 0.4\lambda_L$, two-track decay vertices are selected around the $K_L$ line of flight; the vast majority of the background due to $K_L \rightarrow \pi\mu\nu$, and $\pi^+\pi^-\pi^0$ is rejected by cutting on the $E_{\text{miss}} - P_{\text{miss}}$ distribution, where $P_{\text{miss}}$ and $E_{\text{miss}}$ are the missing momentum and missing energy evaluated in the hypothesis of pion and muon daughter particles. A time of flight technique is used to identify electron and pion tracks. The radiative events are selected by further requiring the detection of a photon with a time of flight compatible with the decay vertex; the cluster position is used to close the kinematics $p_T^2 = 0 = (p_K - p_e - p_\pi - p_\gamma)^2$, and evaluate the energy $E_\gamma$ of the photon. A control sample of $K_L \rightarrow \pi^+\pi^0\pi^0$ decays is used to check the photon efficiency, energy and vertex resolutions. Finally, the signal counts are obtained by fitting the $E_\gamma^*\nu$ and $\theta^*_{e-\gamma}$ bi-dimensional distribution, where $\theta^*_{e-\gamma}$ is the angle between the electron and the photon\(^5\), with signal and background distributions obtained from MC.

The preliminary result for the ratio $R = \text{BR}(K_L \rightarrow \pi e\nu\gamma)/\text{BR}(K_L \rightarrow \pi e\nu(\gamma))$, with the cuts for the exclusive channel $E_\gamma^* > 30$ MeV, and $\theta^*_{e-\gamma} > 20^\circ$, is:

$$R = (0.92 \pm 0.02_{\text{stat}} \pm 0.02_{\text{syst}})\%.$$  

By using the SD spectrum shape evaluated in Ref.\(^7\), we are also able to measure the SD contribution:

$$\text{BR}_{\text{SD}}(K_L \rightarrow \pi e\nu\gamma) = (-3.1 \pm 3.0) \times 10^{-5} \quad \text{and} \quad \text{BR}_{\text{SD}} \leq 2.5 \times 10^{-5} @ 90\% \text{ c.l.},$$

in agreement with theoretical predictions based on chiral perturbation theory calculation\(^7\).

The accuracy of the KLOE result on $R$ is not sufficient to shed light on the discrepancy between previous measurements by KTeV and NA48 collaborations\(^8\)\(^9\). However the analysis of the full KLOE data sample (statistics $\times 5$) will improve the accuracy on both $R$ and $\text{BR}_{\text{SD}}$ results.

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\(^5\)The symbol $\times$ indicates quantities evaluated in the kaon rest frame.
2.4 $K_{L\mu 3}$ form factor slope $\lambda_0$

The knowledge of the $K_L$ scalar form factor $f_0(t)$, where $t$ is the momentum transfer, is relevant for the determination of $V_{us}$ and to test $e/\mu$ universality; typically a linear parametrization is used $f_0(t) \propto 1 + \lambda_0(t/m_{\pi^+}^2)$, where the slope $\lambda_0$ is a parameter to be experimentally determined. A data sample corresponding to $\sim 330$ pb$^{-1}$ has been analyzed; $K_{L\mu 3}$ events are selected requiring a $K_L$ of known momentum and direction, tagged by $K_S \rightarrow \pi^+\pi^-$ decay near the IP. In a fiducial volume extending for $\sim 0.4\lambda_0$, two-track decay vertices are selected around the $K_L$ line of flight; the background due to $K_L \rightarrow \pi \nu \nu$, $\pi^+\pi^-\pi^0$, and $\pi^+\pi^-$ is rejected by cutting on different combinations of $E_{miss}$ and $P_{miss}$ variables, where $E_{miss}$ is evaluated in different masses hypothesis for the daughter particles. The same variables are used to select the signal. A further reduction of the background at the level of $\sim 1.5\%$ is obtained using neural network and time of flight techniques. The analysis of $K_{L\mu 3}$ decays for the measurement of the slope parameter $\lambda_0$ is more complicated than for $K_{Le 3}$ decays$^{11}$ because pure and efficient $\pi - \mu$ separation is much more difficult to achieve. In order to overcome this problem, the analysis aims at measuring $\lambda_0$ through a fit to the distribution of the neutrino energy $E_\nu$, which can be evaluated through a Lorentz transformation of the missing momentum $\vec{P}_{miss}$ in the $K_L$ rest frame. As a consequence the sensitivity on form factor slope $\lambda_0$ is slightly reduced with respect to that achieved from a fit on the $t$ distribution. A combined fit of the neutrino energy spectrum with $K_{L\mu 3}$ results for the vector form factor slopes $\lambda_{+,0}$, $\lambda_{+,1}$ yields the following preliminary result:

$$\lambda_0 = (15.6 \pm 1.8_{\text{stat}} \pm 1.9_{\text{ syst}}) \times 10^{-3}$$

with an accuracy similar to other measurements$^{12}$. The relative statistical accuracy will be in the range $5 - 10\%$ with the analysis of the full KLOE data sample.

3 KLOE summary on $V_{us}$

The most precise test of the unitarity of the CKM matrix can be performed from its first row. Defining $\Delta = |V_{us}|^2 + |V_{us}|^2 + |V_{ub}|^2 - 1$, an accuracy of few parts in $10^{-4}$ on $\Delta$ can be reached. The contribution of $|V_{us}|^2$ is negligible$^{12}$; the determination of $|V_{us}|$ form super-allowed nuclear beta decays gives an uncertainty of $5 \times 10^{-4}$ on $\Delta$.$^{13}$ A similar accuracy can be reached extracting $|V_{us}|$ from the rates of semileptonic kaon decays:

$$\Gamma_i^\mu(K_{e3}(\gamma),\mu3(\gamma)) = \frac{G_F^2 G^{2M^5}}{128\pi^3} |S_{\text{EW}}| |V_{us}|^2 |f_i(0)|^2 I_{e3,\mu3}^i (1 + \delta_{e3,\mu3}^i)$$

(2)

where the index $i$ refers to $K^0 \rightarrow \pi^-$ or $K^+ \rightarrow \pi^0$ transitions for which $G_F^2 = 1$ or $1/2$, respectively, $G$ is the Fermi constant, $M$ is the appropriate kaon mass, and $S_{\text{EW}}$ is a universal short-distance electroweak correction$^{14}$. The $\delta$ term accounts for long-distance radiative corrections depending on the meson charges and lepton masses and, for $K^+$, for isospin-breaking effects. The $f_i(0)$ form factor parametrizes the vector-current transition $K^0 \rightarrow \pi^-$ at zero momentum transfer $t$, while the dependence of vector and scalar form factors on $t$ enters into the determination of the integrals $I_{e3,\mu3}$ of the Dalitz-plot density over the physical region.

The experimental inputs in Eq.(2) are the semileptonic decay widths, evaluated from the $\gamma$-inclusive BR's and from the kaon lifetimes, and the parameters describing the $t$-dependence of the vector and scalar form factors. The KLOE experiment provides measurements for all these quantities$^{c}$ with the only exception of the $K_S$ lifetime; taking the $\tau_S$ from Ref.$^{12}$, the values of $|V_{us}|f_i(0)$ are obtained from KLOE measurements for different decay modes, as listed in Table 1. The best accuracy, $\sim 0.3\%$, is obtained from $K_{Le 3}$ mode, with the error dominated

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$^c$For the recent KLOE measurement of $\tau_\pi$ see Ref.$^{15}$. 

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Table 1: Summary of KLOE measurements of $|V_{us}|f_f(0)$.

| Mode    | $|V_{us}|f_f(0)$ |
|---------|-----------------|
| $K_{L\pi 3}$ | 0.2156(7)       |
| $K_{L\mu 3}$ | 0.2163(10)      |
| $K_{S\pi 3}$ | 0.2154(14)      |
| $K^\pm_3$  | 0.2168(22)      |
| $K^\mp_3$  | 0.2151(30)      |
| **Average** | **0.2158(6)**  |

by the knowledge of $\tau_L$. Using the value of $f_f(0) = 0.961(8)$ evaluated by Leutwyler and Roos\textsuperscript{16}, the value of $V_{us} = 0.2246(20)$ from KLOE average for $K_{\pi 3}$ modes is obtained; using the world average\textsuperscript{13} value of $V_{ud} = 0.97377(27)$, we get $\Delta = (-13 \pm 10) \times 10^{-4}$. A combined analysis of all available experimental results in order to extract $|V_{us}|$ is discussed in detail elsewhere\textsuperscript{17}.

The values of $V_{us}$ obtained from $K_{e 3}$ and $K_{\mu 3}$ decays can also be used to test the universality of $e$ and $\mu$ couplings to the $W$ boson. KLOE results are compatible with a ratio of effective Fermi constants equal to unity:

$$|G(\mu 3)/G(\varepsilon 3)|^2 = 1.0065(98) \quad \text{for } K_L$$

$$|G(\mu 3)/G(\varepsilon 3)|^2 = 0.984(25) \quad \text{for } K^\pm$$

From the measured ratio $\Gamma(K \rightarrow \mu \nu)/\Gamma(\pi \rightarrow \nu \nu)$ of radiation-inclusive kaon and pion widths for $\mu \nu$ decays, the ratio $|V_{us}/V_{ud}|$ can be extracted, using as theoretical inputs the form factor ratio $f_K/f_\pi$ from lattice calculation\textsuperscript{19}, and a radiative correction factor from Ref\textsuperscript{18}. From the precise KLOE measurement\textsuperscript{20} BR$(K^+ \rightarrow \mu^+ \nu) = 0.6366 \pm 0.0009_{\text{stat}} \pm 0.0015_{\text{syst}}$, and using PDG values\textsuperscript{12} for the other experimental inputs, we get:

$$|V_{us}/V_{ud}| = 0.2286 \left(^{+0.02}_{-0.03}\right) .$$

This result can be fit together with the values of $V_{us}$ from KLOE average and $V_{ud}$ from Ref\textsuperscript{23}, yielding the result $V_{us} = 0.232(16)$ and $\Delta = (16 \pm 12) \times 10^{-4}$ with a $\chi^2$ probability of $56\%$, demonstrating the consistency of KLOE measurements, and giving no indication of any violation of CKM unitarity.

4 CPT Test from Bell-Steinberger Relation

The Bell-Steinberger relation (BSR)\textsuperscript{21} relates a possible violation of CPT invariance ($m_{K^0} \neq m_{\bar{K}^0}$ and/or $\Gamma_{K^0} \neq \Gamma_{\bar{K}^0}$) in the time-evolution of the $K^0-\bar{K}^0$ system to the observable $CP$-violating interference of $K_L$ and $K_S$ decays into the same final state $f$. The BSR relation can be written in the following form:

$$\left(\frac{\Gamma_S + \Gamma_L}{\Gamma_S - \Gamma_L} + i \tan \phi_{SW}\right) \left(\frac{R(\varepsilon)}{1 + |\varepsilon|^2} - i \Im(\delta)\right) = \frac{1}{\Gamma_S - \Gamma_L} \sum_f A_S^f(f) A_L^f(f)$$

(3)

without approximations, and phase-convention independent in the exact CPT limit\textsuperscript{22}, where $\varepsilon$ and $\delta$ are the usual $T$ and $CP$ violating parameters in the kaon mixing, respectively, $\phi_{SW}$ is the superweak phase, and $A_{S,L}(f)$ are the decay amplitudes of $K_{S,L}$ into the final state $f$. The solution of Eq. 3 is given by:

$$\begin{pmatrix} \Re(\varepsilon) \\ \Im(\delta) \end{pmatrix} = \frac{1}{N} \begin{pmatrix} 1 + \kappa(1 - 2b) & (1 - \kappa) \tan \phi_{SW} \\ (1 - \kappa) \tan \phi_{SW} & -1 + \kappa \end{pmatrix} \begin{pmatrix} \Sigma_f \Re(\alpha_f) \\ \Sigma_f \Im(\alpha_f) \end{pmatrix},$$

(4)
where on the r.h.s. there are all measurable quantities; the $\alpha_f$ parameters are conveniently defined as follows for the main final states:

$$\alpha_{\pi\pi} \equiv \frac{1}{\Gamma_S} \langle A_L(f), A_S^*(f) \rangle = \eta_f \text{BR}(K_S \to f), \quad f = \pi^0, \pi^- \pi^+ (\gamma),$$ \hspace{1cm} (5)

$$\alpha_{\pi\pi} \equiv \frac{1}{\Gamma_S} \langle A_L(f), A_S^*(f) \rangle = \frac{\tau_{K_S}}{\tau_{K_L}} \eta_f^2 \text{BR}(K_L \to f), \quad f = 3\pi^0, \pi^0 \pi^+ \pi^- (\gamma).$$ \hspace{1cm} (6)

$$\alpha_{\pi\pi} \equiv \frac{1}{\Gamma_S} \sum_{\pi\ell\nu} \langle A_L(\pi\ell\nu), A_S^*(\pi\ell\nu) \rangle + 2\frac{\tau_{K_S}}{\tau_{K_L}} \text{BR}(K_L \to \pi\ell\nu) \Im(\delta)$$

$$= 2\frac{\tau_{K_S}}{\tau_{K_L}} \text{BR}(K_L \to \pi\ell\nu) (A_S + A_L)/4 - \Im(\delta_{\pi\ell\nu})$$ \hspace{1cm} (7)

with $\eta_f = A_L(i)/A_S(i), A_{\pi\pi}$ is the semileptonic charge asymmetry for $K_{SL}, \Im(\delta_{\pi\ell\nu})$ is the \Delta $S - \Delta Q$ violating and CPT conserving parameter for semileptonic decay amplitudes, and $\kappa = \tau_K/K_L, b = \text{BR}(K_L \to \pi\ell\nu), N = (1 + \kappa)^2 + (1 - \kappa)^2 \tan^2 \phi_{SW} - 2b\kappa(1 + \kappa)$.

After having provided all the experimental inputs to Eq.(4) by using KLOE measurements, PDG values, and a combined fit of KLOE and CPLEAR data in order to improve the precision on $\Im(\delta_{\pi\ell\nu})$, we obtain $d$:

$$\Re(\epsilon) = (159.6 \pm 1.3) \times 10^{-5}, \quad \Im(\delta) = (0.4 \pm 2.1) \times 10^{-5}$$

improving by almost a factor three the previous best limit on $\Im(\delta)$ from CPLEAR. \hspace{1cm} (8)

The main limiting factor of the present result is the uncertainty on the phase of $\eta_{\pi^+\pi^-} - \text{parameter entering in} \alpha_{\pi^+\pi^-}$. In fact, using the KLOE upper limit on $\text{BR}(K_S \to 3\pi^0)^{24}$ and the $A_S$ measurement$^{25}, \alpha_{\pi^+\pi^-} \rho_{\pi^+\pi^-}$ and $\alpha_{\pi^+\pi^-}$ do not limit anymore the test sensitivity.

The limits on $\Im(\delta)$ and $\Re(\epsilon)$ can be used to constrain the mass and width difference between $K^0$ and $\bar{K}^0$. Since the total decay widths are dominated by long-distance dynamics, in models where CPT invariance is a pure short-distance phenomenon, it is useful to consider the limit $\Gamma_{K^0} = \Gamma_{\bar{K}^0}$. In this limit we obtain the following bound on the neutral kaon mass difference:

$$-5.3 \times 10^{-10} \text{GeV} < \Delta M < 6.3 \times 10^{-10} \text{GeV} \text{ at } 95\% \text{ CL.}$$

5 Decoherence and CPT Tests Using Kaon Interferometry

The quantum interference between the two kaon decays in the CP violating channel $\phi \to K_SK_L \to \pi^+\pi^-\pi^+\pi^-$ has been observed for the first time by KLOE$^{27}$. A data sample corresponding to $\sim 330 \text{ pb}^{-1}$ has been analysed; the selection of the signal requires two vertices, each with two opposite curvature tracks inside the drift chamber, with an invariant mass and total momentum compatible with the two neutral kaon decays. The experimental resolution on the time difference $\Delta t = |t_1 - t_2|$ in the case of $\pi^+\pi^-$ decays can be improved exploiting the good momentum resolution of the KLOE detector and the closed kinematics of the event. After a kinematic fit, a resolution $\sigma_{\Delta t} \sim 0.9\tau_S$ is obtained. The measured $\Delta t$ distribution can be fitted with the expression:

$$I(\pi^+\pi^-, \pi^+\pi^-; \Delta t) \propto e^{-\tau_{K}\Delta t} + e^{-\tau_{S}\Delta t} - 2(1 - \zeta_{SL})e^{-\frac{\tau_{S+T}}{2}\Delta t} \cos(\Delta m\Delta t)$$

where the quantum mechanical expression in the $\{K_S, K_L\}$ basis has been modified with the introduction of a decoherence parameter $\zeta_{SL}$, and a factor $(1 - \zeta_{SL})$ multiplying the interference

\hspace{1cm} $d$See Ref. $^{33}$ for a detailed discussion on the fit procedure.
term. Analogously, a \( \zeta_{00} \) parameter can be defined in the \( \{K^0, \bar{K}^0\} \) basis\(^8\). After having included resolution and detection efficiency effects, having taken into account the background due to coherent and incoherent \( K_S \)-regeneration on the beam pipe wall, the small contamination of non-resonant \( e^+e^- \rightarrow \pi^+\pi^-\pi^+\pi^- \) events, and keeping fixed in the fit \( \Delta m, \Gamma_S \) and \( \Gamma_L \) to the PDG values, a fit is performed on the \( \Delta t \) distribution, as shown in Fig. 1. The fit results are:

\[
\begin{align*}
\zeta_{SL} &= 0.018 \pm 0.040_{\text{stat}} \pm 0.007_{\text{syst}} \\
\zeta_{00} &= (1.0 \pm 2.1_{\text{stat}} \pm 0.4_{\text{syst}}) \times 10^{-6}
\end{align*}
\]

compatible with the quantum mechanics prediction, i.e. \( \zeta_{SL} = \zeta_{00} = 0 \), and no decoherence effects. In particular the result on \( \zeta_{00} \) has a high accuracy, \( \mathcal{O}(10^{-6}) \), due to the \( CP \) suppression present in the specific decay channel; it improves of five orders of magnitude the previous limit obtained by Berthmann and co-workers\(^8\) in a re-analysis of CPLEAR data. This result can also be compared to a similar one recently obtained in the B meson system\(^{20}\), where an accuracy \( \mathcal{O}(10^{-2}) \) can be reached.

It has been pointed out\(^{30,31}\) that in the context of a hypothetical quantum gravity, \( CPT \) violation effects might occur in correlated neutral kaon states, where the resulting loss of particle-antiparticle identity could induce a breakdown of the correlation of state (1) imposed by Bose statistics. As a result the initial state (1) can be parametrized in general as:

\[
|i\rangle = \frac{1}{\sqrt{2}} \left[ |K^0\rangle|\bar{K}^0\rangle - |\bar{K}^0\rangle|K^0\rangle + \omega \left( |K^0\rangle|\bar{K}^0\rangle + |\bar{K}^0\rangle|K^0\rangle \right) \right],
\]

where \( \omega \) is a complex parameter describing a completely novel \( CPT \) violation phenomenon, not included in previous analyses. Its order of magnitude could be at most

\[
|\omega| \sim \left( \frac{m_K^2}{M_{\text{Planck}}} / \Delta \Gamma \right)^{1/2} \sim 10^{-3}
\]

with \( \Delta \Gamma = \Gamma_S - \Gamma_L \). A similar analysis performed on the same data as before, including in the fit of the \( \Delta t \) distribution the modified initial state Eq.(8), yields the first measurement of the complex parameter \( \omega \)\(^{27}\):

\[
\Re(\omega) = \left( 1.1^{+8.7}_{-5.3} \pm 0.9 \right) \times 10^{-4} \quad \Im(\omega) = \left( 3.4^{+4.8}_{-5.0} \pm 0.6 \right) \times 10^{-4} ;
\]

with an accuracy that already reaches the interesting Planck's scale region. Other interesting results related to possible decoherence and \( CPT \) violation in the quantum gravity framework are discussed in Ref.\(^{27}\).
Conclusions

New preliminary results on $K_S \rightarrow e^+e^-$, $K_S \rightarrow \gamma\gamma$, $K_L \rightarrow \pi e\nu\gamma$, and $K_{L\beta3}$ form factor have been presented, and some recent results on $V_{us}$, CPT invariance and QM tests reviewed. The analysis of the full data sample of about 2.5 fb$^{-1}$ is in progress, and new or improved results will be available in the next future. In the meanwhile the possibility to continue the KLOE physics program at DAΦNE with an improved luminosity at the $\phi$ peak up to $\sim 10^{33}\text{cm}^{-2}\text{s}^{-1}$ is strongly considered.

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IV - Astrophysics-and Cosmology-related issues - dark matter, axions...
Photons as a Probe of Minicharged Particles

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Low energy experiments with photons can provide deep insights into fundamental physics. In this note we concentrate on minicharged particles. We discuss how they can arise in extensions of the standard model and how we can search for them using a variety of laboratory experiments.

1 Introduction – Light particles coupled to photons

Light particles weakly coupled to photons appear in a variety of extensions of the standard model. A prominent example is the axion invented to solve the strong CP problem. The axion is an example of a pseudo-scalar particle coupled to two photons via a dimension five interaction,

\[ \mathcal{L}_{\text{int}}^{(-)} = -\frac{1}{4}g\phi^{(-)}F_{\mu\nu}F^{\mu\nu} = g\phi^{(-)}(E \cdot B), \]

(1)

where the coupling constant \( g \) has dimensions 1/Mass. Other examples of such light spin-0 bosons are familons, Majorons, the dilaton, and moduli, to name just a few. Independent of their origin light particles coupled to two photons as in Eq. (1) are often called axion-like particles or ALPs.

ALPs can be constrained by a variety of astrophysical observations. However, these bounds can be avoided in more complicated models making it desirable to have clean and controlled laboratory tests.

Two types of experiments are particularly noteworthy. First there are experiments that look for changes in the polarization when a laser beam passes through a strong magnetic field as depicted in Fig. 1. This is a disappearance experiment where the produced particles are not detected. A pioneering experiment of this type was done by the BFRT collaboration and produced limits on the allowed couplings and masses. Recently, the PVLAS experiment reported the observation of a non-vanishing rotation signal. This has sparked a significant amount of theoretical work as well as planning and construction of new experiments. The Q&A experiment has already published some data but its sensitivity is not yet sufficient to test PVLAS.

Second, there are photon regeneration experiments or “light shining through walls” experiments, as shown in Fig. 2, where the produced particles are reconverted into photons and then detected. BFRT also run a setup of this type. And, particularly interesting, most upcoming experiments will be of this type (or at least have a “light shining through walls” stage).

*For a scalar particle the \( F \) in Eq. (1) has to be replaced by an \( F \).
ALPs have zero electric charge. What about light charged particles? At first one might be tempted to exclude this possibility simply by saying if it is charged and lighter than an electron it is excluded by a huge number of experiments. However, this is implicitly based on the assumption that charge is quantized and the smallest quantum is not much smaller than the charge of the electron. Although strong bounds on the charges of neutrons, atoms and molecules suggest the idea that charge quantization is a fundamental principle one needs physics beyond the standard model to enforce charge quantization. One possibility would be the existence of magnetic monopoles as demonstrated by Dirac’s seminal argument. However, many extensions of the standard model do indeed contain particles with small electric charges.

The interaction for such particles is the standard minimal coupling but with a small fraction $\varepsilon$ of a unit electric charge. For example for Dirac fermions it reads,

$$\mathcal{L}_{\text{int}}^{\text{D}} = \varepsilon e \bar{\psi}_c \gamma_{\mu} \psi_\ell A^\mu.$$  \hspace{1cm} (2)

As we will see in Sect. 3 such an interaction can be tested in optical experiments (cf. Fig. 4) as well as in experiments with strong electric fields as shown in Fig. 5.

2 Minicharged particles in paraphoton models

Minicharged particles arise most naturally in models with extra U(1) gauge degrees of freedom so called paraphotons. In this section we briefly review how kinetic mixing leads to minicharged particles.

Let us look at the simplest model with two U(1) gauge groups, one being our electromagnetic U(1), the other a hidden sector U(1) under which all Standard Model particles have zero charge. The most general Lagrangian allowed by the symmetries is,

$$\mathcal{L} = -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} - \frac{1}{2} \chi F_{\mu\nu} B^{\mu\nu},$$ \hspace{1cm} (3)
where $F^{\mu\nu}$ is the field strength tensor for the ordinary electromagnetic U(1) gauge field $A^\mu$ and $B^{\mu\nu}$ is the field strength for the hidden sector U(1) field $B^\mu$, i.e. the paraphoton. The first two terms are the standard kinetic terms for the photon and paraphoton fields, respectively. Because the field strength itself is gauge invariant for U(1) gauge fields the third term is also allowed by the gauge symmetries (and Lorentz symmetry). This term corresponds to a non-diagonal kinetic term, i.e. a so called kinetic mixing.

The kinetic term can be diagonalized by a shift

$$B^\mu \rightarrow B^\mu - \chi A^\mu.$$  \hspace{1cm} (4)

Aside from a multiplicative renormalization of the gauge coupling $e^2 \rightarrow e^2 / (1 + \chi^2)$ the visible sector fields remain unaffected by this shift.

Let us now assume that we have a hidden sector fermion $f$ that has charge one under $B^\mu$. Applying the shift Eq. (4) to the coupling term we find,

$$e_h \bar{f} B f \rightarrow e_h \bar{f} B f - \chi e_h \bar{f} A f,$$  \hspace{1cm} (5)

where $e_h$ is the hidden sector gauge coupling. We can read off that the hidden sector particle now has a charge

$$ee = -\chi e_h,$$  \hspace{1cm} (6)

under the visible electromagnetic gauge field $A^\mu$ which has gauge coupling $e$. Since $\chi$ is an arbitrary number the fractional electric charge $e$ of the hidden sector fermion $f$ is not necessarily integer.

For small $\chi \ll 1$

$$|e| \ll 1$$  \hspace{1cm} (7)

and $f$ becomes a minicharged particle. From now on we will concentrate on this case.

To conclude this section let us comment on the origin of the kinetic mixing term in Eq. (3) (for more details see, e.g., 49,50,51,24). First of all it should be stressed that the kinetic mixing term is allowed by all symmetries and therefore a free parameter from the viewpoint of an effective low energy field theory. Having said this it is also clear that such a term will typically be generated by loop diagrams in quantum field theory. For example if we have a heavy particle that is charged under both the electromagnetic as well as the hidden sector U(1) gauge group we find a diagram as in Fig. 3(a) which automatically generates a kinetic mixing term. In string theory a similar diagram with an open string going around the loop exists (Fig. 3(b)). In D-brane models of string theory stacks of D-branes generate U(N) gauge groups. The diagram Fig. 3(b) can then be understood as a closed string exchange between two stacks of D-branes. One may imagine that we live on a stack of such D-branes, the "visible" sector, that communicates via such a closed string with another stack of D-branes, the "hidden" sector (cf. Fig. 3(c)). In this way observing kinetic mixing is a first step towards observing the hidden sector which is a common feature in string theory models.

3 Searching minicharged particles in the laboratory

Let us now look how we can actually search for these minicharged particles in the laboratory.

A classic probe of minicharged particles in the laboratory is the invisible decay of orthopositronium. The current limit from this type of experiments is $e < 3.4 \times 10^{-5}$.

\footnote{Light particles with charge $e = O(1)$ are excluded by experiments and very massive particles give negligible contributions in experiments such as PVLAS or the upcoming optical experiments.}

\footnote{As for ALPs the astrophysical bounds are quite strong, $e \lesssim 10^{-14}$ (see, e.g., 53,54,55,56,57), but may be circumvented in some models (note, however, 28 for a bound that may be more difficult to avoid).}
Figure 3. (a) One-loop diagram which contributes to kinetic-mixing in field theory, (b) its equivalent in open string theory and (c) reinterpretation of (b) as a closed string exchange in the context of D-brane models.

Figure 4: Rotation (left) and ellipticity (right) caused by the existence of a light minicharged particle. In a homogeneous magnetic background $\mathcal{B}$, the interaction Eq. (2) can convert the laser photons into pairs of charged particles (left). The conversion probability depends on the relative orientation of the magnetic field and the laser polarization, resulting in an overall rotation. In a similar manner virtual pair production (right) in the magnetic field leads to an orientation dependent index of refraction. This causes a phase shift that appears as an ellipticity in the outgoing beam.

In the small mass range optical experiments provide an even more powerful tool in the search for minicharged particles. Again we can test for changes in the polarization of a laser beam after it has passed through a strong magnetic field. As for ALPs we can have real and virtual production of particles but now it is pair production instead of single particle production. The relevant processes are depicted in Fig. 4.

For small masses data from the BFRT and Q&A experiments constrain

$$\epsilon < 1.2 \times 10^{-6}, \quad \text{for} \quad m_c \lesssim 10^{-2} \text{eV}. \quad (8)$$

Interpreted as a minicharged particle effect the observed rotation in the PVLAS experiment would suggest particles with a mass $m_c \lesssim 0.1 \text{eV}$ and a charge in the range $(0.7 - 1.2) \times 10^{-6}$. This makes this interpretation testable in the immediate future (see also below).

Another way to search for minicharged particles is to employ Schwinger pair production in strong electric fields. In a strong electric field charged particles gain energy when they are separated along the lines of the electric field. Now, if a virtual particle-antiparticle pair generated by a vacuum fluctuation is separated by a large enough distance such that the energy gain in the electric field is bigger than the rest mass of the particles the virtual pair becomes real. In other words a strong electric field can “decay” into particle-antiparticle pairs of charged particles. This process is similar to tunneling. An energy barrier (rest mass) with finite extent (distance between the particles that is sufficient such that the energy gain in the electric field can compensate for the rest mass) can be quantum mechanically crossed. As expected the rate is exponentially small if the barrier is high (large mass) and the distance is large (high mass, weak electric field and small charge). However, if the mass is small pair production can be quite effective.
Figure 5: "Dark current flowing through the walls" experiment (left). In a cavity with strong electric fields, Schwinger pair production produces pairs of minicharged particles. The produced particles have typically momenta along the lines of the electric field. Positive particles flying in one direction and negative particles in the opposite direction. This results in a current. This current is, however, made up of very weakly interacting particles and can therefore pass through thick layers of material (in contrast to, e.g., a current made up from electrons). This current can then be measured on the other side of the wall (e.g. by measuring the generated magnetic field). For sufficiently strong electric field pair production is very efficient and the currents can reach values measurable with current technology. The plot in the right panel shows a (very optimistic) estimate for a current generated in a cavity of length 20 cm and radius 10 cm with a field strength of \( \sim 15 \text{ MV/m} \). From top to bottom the lines correspond to currents from 1 \( \mu \text{A} \) to 1 \( n \text{A} \).

One possibility to generate such strong electric fields is to use accelerator cavities.

The remaining question is how do we know that we have produced minicharged particles in the cavity. Direct detection seems difficult because the cross section decreases with \( \sim \epsilon^2 \) and becomes quite small for small \( \epsilon \). Here, it is useful that once Schwinger pair production sets in, it really can produce a lot of particles. Such a massive production of particles drains a macroscopic amount of energy from the cavity which can be detected (e.g. it would lead to a decrease in the quality factor of the cavity). From available data on the energy loss of cavities for the TESLA accelerator one can infer\(^23\) strong limits on the existence of light minicharged particles \( \epsilon \lesssim 10^{-6} \) for masses \( m_e \lesssim 10^{-3.5} \text{ eV} \).

Using Schwinger pair production process one could even set up a detection experiment (measuring the energy loss is a disappearance experiment) by measuring currents generated in the cavities as depicted in Fig. 5. With available technology such an experiment has chances to probe the interesting region of \( \epsilon \sim \text{few} \times 10^{-7} \) favored by a minicharged particle interpretation of the PVLAS data.

4 Conclusions and outlook

Light particles coupled to photons appear in a wide variety of possible extensions of the standard model. In particular minicharged particles can arise from models with extra \( U(1) \) gauge degrees of freedom. Such particles can be searched for in experiments with photons as, e.g., in experiments that shine laser light through strong magnetic fields or by searching for Schwinger pair production in strong electric fields. Another promising approach is to search for the effects of the additional \( U(1) \) degrees of freedom\(^22\).

Inspired by the PVLAS observation several new experiments suitable for the search for light particles coupled to photons are in planning or are already under construction. Additional experiments such as the "Dark current flowing through a wall" could be built with present technology. These experiments will not only allow to test the PVLAS result but will probe whole classes of viable extensions of the standard model.
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Could a $\gamma$ Line Betray the Mass of Light Dark Matter?

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We\textsuperscript{1} compute the pair annihilation cross section of light dark matter scalar particles into two photons, and discuss the detectability of the monochromatic line associated with these annihilations.

1 Introduction

The need for Cold Dark Matter (DM) to describe and understand how structures in the universe hold together, has become increasingly pressing with the impressive amount of observational data collected in the last ten years. Alternatives to DM, like modifications of gravity, are being put to critical and maybe fatal test by recording maps of gravitational lensing. Indeed, the separation recently observed in colliding clusters between the (maybe modified) gravitational deviation of light and the normal matter that causes it, seems very contrived with modified gravity, and very natural if collisionless DM is the main source of gravity. Questions about the nature of DM and its non-gravitational detection are therefore more relevant than ever.

In this context, the precise determination by INTEGRAL/SPI\textsuperscript{2} of the characteristics of the 511 keV line emitted in our galaxy is intriguing. Indeed, it implies without any doubt that the central bulge of our galaxy is a strong source of positrons. Astrophysical sources like Low Mass X-ray Binaries (LMXB) and Type 1a Supernovae (SN1A)\textsuperscript{3} cannot naturally explain why this source is at the same time steady, extended and absent in the disk. On the contrary, production of positrons through DM annihilation is naturally steady and concentrated in the bulge where the DM density increases: a fit of the needed DM density profile can even be attempted\textsuperscript{4}, yielding a reasonable NFW profile $\rho(r) \sim r^{-1}$ for DM annihilation at rest, in opposition to a less reasonable $\rho \sim r^{-2}$ for DM decay.
In order to maximize the electron-positron annihilation channel, such DM must be light (LDM) \(^ {5, 6} \), at least below muon pair threshold: \( m_{dm} < 100 \text{ MeV} \). More constraining upper bounds can be obtained by a careful study of final state radiation processes (dm dm\(^*\) \rightarrow e^+ e^- \gamma \gamma) and positron annihilation in flight, both producing continuous gamma ray spectra which increase with \( m_{dm} \). From the first, \( m_{dm} < 35 \text{ MeV} \) is obtained \(^7 \), and \( m_{dm} < 20 \text{ MeV} \) from comparing the second with error bars on the measured spectra. On the other hand, \( m_{dm} \) should be higher than 2 MeV to avoid spoiling nucleosynthesis \(^8 \), and higher than 10 MeV if there is a significant coupling to neutrinos which can alter supernova explosions \(^9 \), but such is not necessarily the case.

A fairly unique viable model satisfying all the above constraints contains scalar DM particles with \( m_{dm} \sim 10 \text{ MeV} \), annihilating at rest in the galactic bulge into \( e^+ e^- \) pairs via \( t \)-channel exchange of heavy (\( > 100 \text{ GeV} \)) fermions \( F_e \). Given the large local DM abundance inferred from the rotation curve, the annihilation cross-section yielding the observed positron source is however too small to explain a correct relic density inferred e.g. from cosmic microwave background measurements. A further light vector particle can then be invoked to mediate an \( s \)-channel annihilation process: being velocity dependent, this process becomes dominant in the early universe and can independently be adjusted to the relic density.

If correct, such a model \(^5 \) would profoundly alter the road to unification in particle physics. It therefore seems important to look for other experimental cross-checks. The simplest and most convincing one would be the discovery of another gamma ray line, from the process \( dm \ dm\(^*\) \rightarrow \gamma \gamma \). In the following, we show \(^4 \) this line is unavoidable in such a model, estimate its intensity, and discuss its observability.

2 Dark matter annihilation cross section into two photons

The model considered is specified by the Lagrangian \( \mathcal{L} = \bar{\psi}_F (c_F P_R + c_{\bar{F}} P_L) \psi_x \phi_{dm} + h.c \) where \( P_{R,L} \) are the chiral projectors \((1 \pm \gamma_5)/2 \). The relevant annihilation diagrams are box-diagrams containing 1, 2 or 3 heavy fermions \( F_e \). Assuming that \( dm \neq dm^* \) (which fixes the circulation of arrows), there are 6 diagrams, taking into account permutation of the 2 photon external legs.

From naive power counting, each box is logarithmically divergent. However, gauge invariance dictates a result proportional to \( F^2_{\mu\nu} \) rather than \( A^2_{\mu\nu} \). This requires 2 powers of external momenta, so that the integrand must in fact converge like \( d^4 k/k^3 \) for large loop momenta \( k \). In the limit \( m_{F_e} \gg m_{e, dm} \) (relevant due to LEP and other collider/accelerator constraints), the contribution of momenta larger than \( m_{F_e} \) is \( \sim 1/m_{F_e}^2 \). The leading \( 1/m_{F_e} \) term can thus be safely obtained by expanding the integrand in powers of \( 1/m_{F_e} \) and keeping only the first term.

\[
\begin{align*}
\frac{dm}{dm^*} & \quad \frac{e_{R,L}}{e_{L,R}} \quad F \quad e \quad e \\
\frac{dm}{dm^*} & \quad \frac{e_{R,L}}{e_{L,R}} \quad F \quad e \quad e \\
\frac{dm}{dm^*} & \quad \frac{e_{R,L}}{e_{L,R}} \quad F \quad e \quad e \\
\frac{dm}{dm^*} & \quad O(\frac{1}{m_{F_e}^2}) \quad \L_{eff} = \frac{1}{m_{F_e}^2} \psi_x \phi_{dm} \phi_{dm} \bar{\psi}_x (a + ib\gamma_5) \psi_x
\end{align*}
\]
This corresponds to “pinching” the box with one $F_e$ into a triangle involving only electrons and an effective $d_m-d_m-e-e$ coupling given by the real couplings $a, b$ given by $a + ib = c^2 F_e$. For this set-up, computing the cross-section is a loop-textbook exercise for which we find:

$$
\sigma_{\gamma \gamma} v_r = \frac{\alpha^2}{(2\pi)^3 m_{F_e}^2 m_{d_m}^2} \times \left[ b^2 \left| 2C_0 m_{d_m}^2 \right|^2 + a^2 \left| 1 + 2C_0 (m_{e}^2 - m_{d_m}^2) \right|^2 \right].
$$

$C_0$ is a function of $m_e$ and $m_{d_m}$ given by the Passarino-Veltman scalar integral. For $m_{d_m} > m_e$, this function develops an imaginary part corresponding to the formation of a real $e^+e^-$ pair subsequently annihilating into 2 photons, and giving the largest contribution for masses above 1 MeV.

For $m_{d_m} \ll m_e$, $C_0$ behaves as $-1/(2m_e^2) + m_{d_m}^2/(3m_e^4)$, so that both terms of the cross section behave as $m_{d_m}^2/(m_e m_{F_e})^2$. This limit is relevant to estimate the effect of heavier particles than the electron in the loop. For example, the contribution of the $\tau$ lepton could be significant if the corresponding couplings $(a_{\tau}, b_{\tau})$ are larger than $\approx (m_{\tau}/m_{d_m}) \times (a_e, b_e) \times (m_{F_e}/m_{F_e})$ (with $m_{d_m} < m_{\tau}$), i.e. if they scale at least like usual Yukawa couplings. Since an independent detailed analysis is required to check whether or not such couplings can pass particle physics constraints, we prefer giving a conservative estimate based on the electron contribution only. The latter cannot be turned off without losing the 511 keV line signal. It therefore constitutes a safe lower bound for assessing the detectability of the line at $E_{\gamma} = m_{d_m}$.

Within the pinch approximation, the cross-section relevant for the origin of the 511 keV emission is:

$$
\sigma_{511} v_r = \frac{\beta_e}{4\pi m_{F_e}^2} \left( a^2 \beta_e^2 + b^2 \right)
$$

with $\beta_e = \sqrt{1 - m_e^2/m_{d_m}^2}$, which indeed for $b = 0$ reduces to the expression used for large $m_{F_e}$. After careful comparison with SPI data, Ascasibar et al. found

$$
\sigma_{511} v_r = 2.6 \times 10^{-30} \left( \frac{m_{d_m}}{\text{MeV}} \right)^2 \text{cm}^3/\text{s}.
$$

The $\gamma \gamma$ annihilation cross-section is then also determined by this measurement in terms of the ratio of annihilation branching ratios:

$$
\eta \equiv \frac{\sigma_{\gamma \gamma}}{\sigma_{511}} = \frac{a^2}{2\pi^2 \beta_e^2} \frac{m_{d_m}^2}{m_{d_m}^2} \frac{a^2}{a^2 + b^2} \left[ 1 + 2 \left( m_e^2 - m_{d_m}^2 \right) \right] C_0^2 + b^2 \left| 2C_0 (m_{e}^2 - m_{d_m}^2) \right|^2.
$$

(1)

As announced, this ratio cannot vanish, whatever the value of $a/b$, so that a minimum $\gamma \gamma$ flux is guaranteed. As $m_{d_m}$ approaches $m_e$ from above, the ratio increases like $\beta_e^{-3}$ for a pure scalar coupling ($b = 0$) and like $\beta_e^{-1}$ for an axial one ($a = 0$). The ratio decreases almost linearly with the dark matter mass for $m_{d_m} > 1$ MeV. In the table below, we give typical values of the ratio $\eta$ for the most conservative case (i.e. $a = 0$, $\beta_e^{-1}$):

<table>
<thead>
<tr>
<th>$m_{d_m}$(MeV)</th>
<th>0.52</th>
<th>1</th>
<th>5</th>
<th>20</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\eta (a = 0)$</td>
<td>8.8 $10^{-6}$</td>
<td>1.4 $10^{-5}$</td>
<td>3.6 $10^{-6}$</td>
<td>8.1 $10^{-7}$</td>
</tr>
</tbody>
</table>

Notice that the simple guess applied to the case of decaying DM

$$
\eta_{\text{guess}} \approx \frac{a^2 m_{d_m}^2}{2\pi^2 \beta_e^2}
$$


Figure 1: Left: flux of the monochromatic $E_\gamma = m_{\text{dim}}$ line from a 8 degree cone around the galactic center. Right: significance of the monochromatic $E_\gamma = m_{\text{dm}}$ line above the continuum background for one year of observation with an ideal detector of 1 m$^2$ and a 10$^{-3}$ energy resolution.

increases instead of decreasing with $m_{\text{dm}}$. For a typical mass of 10 MeV, this guess overestimates the monochromatic flux by a factor 635 with respect to our result (Eq.1). As we will see in the next section, such a factor is crucial to the line observability.

3 Detectability of the monochromatic line

A few experiments have already scanned the energy range above the electron mass. The instruments on board of INTEGRAL for example have been designed to survey point-like objects as well as extended sources over an energy range between 15 keV-10 MeV. The instrument INTEGRAL/SPI itself is a spectrometer designed to monitor the 20 keV-8 MeV range with excellent energy resolution. Therefore a legitimate question is whether or not the line $E_\gamma = m_{\text{dm}}$ could have been (or could be) detected by the same instrument that has unveiled the 511 keV signal. This essentially depends on the ratio $\eta$ as given above, and on the background.

The 511 keV emission has been measured with a $\sim 10\%$ precision\textsuperscript{11} to be

$$\langle I_{511} \rangle = 6.62 \times 10^{-3} \text{ ph cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$$

inside a region that extends over $350^o < l < 10^o$ in longitude and $|b| < 10^o$ in latitude. If this emission originates from a NFW distribution of LDM species around the galactic center with a characteristic halo radius of 16.7 kpc, the signal from the inner 5$^o$ is found\textsuperscript{4} to be

$$\langle I_{511}(5^o) \rangle = 1.8 \times 10^{-2} \text{ ph cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$$

once the SPI response function is taken into account and the instrumental background is properly modeled. If the positron propagation is negligible, then the map of the 511 keV emission should correspond to that of the LDM annihilations.
Within this approximation, we expect the spatial distributions of both the 511 keV and the two-gamma ray lines to be identical and their intensities to be related by the ratio $\eta$:

$$\langle I_{\gamma\gamma}(\theta_{\gamma\gamma}) \rangle = \frac{\theta_{511}}{\theta_{\gamma\gamma}} \frac{\langle I_{511}(\theta_{511}) \rangle}{(1 - 3f/4)}.$$ 

This expression is approximately valid as long as the angular radii $\theta_{\gamma\gamma}$ and $\theta_{511}$ of the regions monitored by the gamma-ray spectrometer are small. In what follows, the fraction $f$ of positrons forming positronium has been taken equal to 93%. This is in perfect agreement with the positronium fraction later derived, i.e. $f_{Ps} = 0.92 \pm 0.09$. The monochromatic line flux

$$\phi_{511}(\theta_{\gamma\gamma}) = \pi \theta_{\gamma\gamma}^2 \langle I_{511}(\theta_{\gamma\gamma}) \rangle$$

has been plotted in Fig. 1 in the case of the LDM model with $F$ exchange and assuming a NFW profile. The angular radius $\theta_{\gamma\gamma} = 8^\circ$ corresponds to the field of view of the satellite. For typical LDM masses in the MeV range, the expected flux is about three orders of magnitude below the claimed INTEGRAL/SPI line sensitivity (which is about $2.5 \times 10^{-5}$ ph cm$^{-2}$ s$^{-1}$ after $10^9$ seconds). An unrealistic exposure of 30,000 years would thus be required in order to detect the $E = m_{dm}$ line. When the LDM species is degenerate in mass with the electron, the flux is only a factor of 25 below the SPI detection limit (assuming a pure scalar coupling $b=0$). As long as the mass difference $m_{dm} - m_e$ does not exceed 0.1 MeV, it is roughly comparable with the expected 478 keV line signal emitted by Novae, that is about $10^{-7}$ ph cm$^{-2}$ s$^{-1}$.

SPI sensitivity is limited by the instrumental background that arises mostly from cosmic rays impinging on the apparatus and activating the BGO scintillator. On the contrary, the absolute sensitivity of an ideal instrument is purely limited by the gamma-ray continuum background. This emission has been recently estimated

$$I_{BG}(E) = 1.15 \times 10^{-2} E^{-1.82} \text{ph cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{MeV}^{-1},$$

inside the central region that extends over $350^\circ < l < 10^\circ$ in longitude and $|b| < 10^\circ$ in latitude. The energy $E$ is expressed in units of MeV.

We thus estimate the significance $\Sigma \equiv \text{signal/\sqrt{background}}$ for the LDM line to emerge above this background (assuming it is isotropic) to be

$$\Sigma = \sqrt{\pi} \theta_{511} \frac{\langle I_{511}(\theta_{511}) \rangle}{(1 - 3f/4)} \eta \frac{\sqrt{S_0 T_0}}{I_{BG} \Delta E_0},$$

with $S_0$ the surface of the detector, $T_0$ the exposure time, $I_{BG}$ the above-mentioned continuum background intensity and $\Delta E_0$ the energy resolution. The significance $\Sigma$ (displayed as a function of the dark particle mass in Fig. 1 for a surface of 1 m$^2$, an exposure duration of $T_0 = 1$ year and an energy resolution of 0.1%) indicates that those values would theoretically allow to extract the minimal guaranteed signal computed at 3 standard deviations above background for all relevant LDM masses below 30 MeV.

There is nothing to be gained by narrowing the angular aperture $\theta_{\gamma\gamma}$ because, for the assumed NFW profile, the signal increases linearly with this angular radius, as does the square root of an isotropic background.

In contrast, note that the monochromatic line should be extremely narrow: its width is expected to be about a few eV which experimentally is very challenging if one compares
it with the present SPI sensitivity that is about $10^{-3}$ at MeV energies. At lower energies, there are nevertheless instruments, e.g. X-ray CCD, bolometers, Bragg spectrometers which are able to resolve eV widths. A significant improvement on the resolution $\Delta E_0$ at higher energies would probably be necessary in order to reach a large enough significance and ensure detection. Indeed, an effective surface of 1m$^2$ might be hard to attain in space.

Next generation instruments such as AGILE/(super AGILE) or GLAST, which in principle could be more promising, will probably be limited by the energy range that they are able to investigate. Future instruments might nevertheless be able to see this line if their energy resolution and sensitivity are improved by a large factor with respect to SPI present characteristics.

Maybe a better chance to detect this line is to do observations at a high latitude and a longitude slightly off the galactic centre. In this case, indeed, the background should drop significantly (the density of dark clouds has been measured recently\textsuperscript{15}) but the line flux may decrease by a smaller factor.

**Acknowledgement**

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13. J. P. Roques et al. 1000.
Dark Matter from the Inert Doublet Model

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The Inert Doublet Model is an extension of the Standard Model including one extra “Inert scalar doublet” and an exact $Z_2$ symmetry. The “Inert scalar” provides a new candidate for dark matter. We present a systematic analysis of the dark matter abundance assuming the standard freeze-out mechanism and investigate the potentialities for direct and gamma indirect detection. We show that the dark matter candidate saturates the WMAP dark matter density in two rather separate mass ranges, one between 40 and 80 GeV, the other one over 400 GeV. We also show that the model should be within the range of future experiments, like GLAST and EDELWEISS II or ZEPLIN.

1 Introduction

Many evidences for the existence in the universe of dark matter has been put forward over the years. It can be inferred from the dynamics of galaxies and of clusters of galaxies, from analysis of the CMBR, from structure formation, etc. The question that arises then is what is the nature of dark matter and which extension of the Standard Model do we have to consider in order to account for these observations? A profusion of models of dark particles have been proposed over the years and it is much hoped that present and forthcoming experiments will throw some light on the matter (for a review, see for instance 1).

In these proceedings, we study a simple extension of the Standard Model with one extra scalar doublet and an exact $Z_2$ symmetry. In this framework, the candidate for dark matter is one of the two neutral scalars arising from the extra doublet. The latter was called “Inert doublet” by Barbieri et al in 2 because it has no direct coupling to matter fields. However it couples to the standard gauge fields. The phenomenology of its neutral and charged components is quite simple and yet very rich.

This is not the only attractive feature of the model. As was pointed out in 2, the Inert Doublet Model (IDM) could allow for a Higgs mass up to 500 GeV still fulfilling the LEP Electroweak...
2 Short description of the Model

The IDM is a particular two Higgs doublet model, in which one of the doublet, \( H_1 \), plays the role of the standard Brout-Englert-Higgs doublet while the second one, \( H_2 \), is the source for dark matter candidates. In order to guarantee the stability of the dark matter particles, one invoke a \( Z_2 \) symmetry under which all Standard Model fields are even and

\[
H_1 \rightarrow H_1 \quad \text{and} \quad H_2 \rightarrow -H_2.
\]

This discrete symmetry also prevents the appearance of flavor changing neutral currents in this model. Moreover, we assume that \( Z_2 \) is not spontaneously broken. This model was first introduced by Deshpande and Ma\(^8\) (see also\(^7\)), and the dark matter aspect was recently discussed by Cirelli et al\(^8\) and Barbieri et al\(^2\). Their initial purpose and some of their assumptions were nevertheless not exactly identical. In addition, the neutral scalar reaching the dark matter WMAP abundance was found to be in the mass range of 60 to 75 GeV for Barbieri et al\(^2\) while for Cirelli et al\(^8\) it was of order of 430 GeV. We first study the details of the model before to elucidate this apparent incompatibility.

The most general potential of the model can be written as

\[
V = \mu_1^2 |H_1|^2 + \mu_2^2 |H_2|^2 + \lambda_1 |H_1|^4 + \lambda_2 |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 + \lambda_4 |H_1|^2 |H_2|^2 + \frac{\lambda_5}{2} \left[ (H_1^* H_2) + (H_2^* H_1) + h.c. \right] (1)
\]

The vacuum expectation value of \( H_1 \) is given by \( \langle H_1 \rangle = \frac{\nu}{\sqrt{2}} \) with \( \nu = \sqrt{-\mu_1^2 / \lambda_1} = 248 \text{ GeV} \), while assuming for simplicity \( \mu_2^2 > 0 \), we have \( \langle H_2 \rangle = 0 \). The mass of the Higgs particle \( h \) is \( M_h^2 = -2\mu_1^2 \equiv 2\lambda_1 \nu^2 \) while the mass of the charged, \( H^+ \), and two neutral, \( H_0 \) and \( A_0 \), components of the field \( H_2 \) are given by

\[
M_{H^+}^2 = \mu_2^2 + \lambda_5 v^2 / 2 \\
M_{H^0}^2 = \mu_3^2 + (\lambda_3 + \lambda_4 + \lambda_5) v^2 / 2 \\
M_{A_0}^2 = \mu_3^2 + (\lambda_3 + \lambda_4 - \lambda_5) v^2 / 2. (2)
\]

For appropriate quartic couplings, \( H_0 \) or \( A_0 \) is the lightest component of the \( H_2 \) doublet. In the absence of any other lighter \( Z_2 \) odd field, either one is a candidate for dark matter. For definiteness we choose \( H_0 \). All our conclusions are unchanged if the dark matter candidate is \( A_0 \) instead. Following\(^2\), we parameterize the contribution from symmetry breaking to the mass of \( H_0 \) by \( \lambda_L = (\lambda_3 + \lambda_4 + \lambda_5) / 2 \), which is also the coupling constant between the Higgs field \( h \) and our dark matter candidate \( H_0 \).

3 Dark matter abundance

As in\(^2\) and in\(^8\), we consider a thermal production of the cold relic \( H_0 \). We have computed the relic abundance of \( H_0 \) using micrOMEGAs2.0, a new and versatile package for the numerical calculation of dark matter abundance from thermal freeze-out\(^9\).

We first present the results for fixed Inert scalars mass differences in the contour plots of figure 1. We work in the \((M_{H_0}, \mu_2)\) plane, as a result the diagonal line corresponds to \( \lambda_L = 0 \), i.e. to no coupling between \( H_0 \) and the Higgs boson. Away from this line, \( \lambda_L \) increases, with
\[ \lambda_L < 0 \text{ (resp. } \lambda_L > 0) \text{ above (resp. below) the diagonal. Also, we write } \Delta M A_0 = M_{A_0} - M_{H_0} \text{ and } \Delta M H_c = M_{H^+} - M_{H_0}. \]

The shaded areas in the plots correspond to regions that are excluded by several constraints. In order not to conflict with LEP data, the mass of the \( H^+ \) should be larger than 79.3 GeV and \( M_{H^0} + M_{A^0} \leq M_{H^0} \). These constraints translate into the excluded region 1 on the plot on the left of figure 1. The vacuum stability constraint contributes largely to the exclusion of \( \lambda_L < 0 \) couplings. This corresponds to shaded area in the domains \( \mu_2 > M_{H_0} \) in the plots of figure 1. The remaining shaded regions are excluded due to large couplings \( |\lambda_i| > 4\pi \). Moreover regions where the couplings range as \( 1 < |\lambda_i| < 4\pi \), which is still tolerable, are shown with horizontal lines. The areas between two dark lines correspond to regions of the parameter space such that \( 0.094 < \Omega_{DM} h^2 < 0.129 \), the range of dark matter energy densities consistent with WMAP data.

We immediately see that there are two qualitatively distinct regimes, depending on whether the \( H_0 \) is lighter than the \( W \) and \( Z \) and/or the Higgs boson. For the low mass regime, let us study figure 1, the plot on the left. The two processes relevant below the \( W \), \( Z \) or \( h \) threshold are the \( H_0 \) annihilation through the Higgs and \( H_0 \) coannihilation with \( A_0 \) through \( Z \) exchange. Both give fermion-antifermion pairs, the former predominantly into \( b\bar{b} \). Coannihilation into a \( Z \) may occur provided \( \Delta M A_0 \) is not too important, roughly \( \Delta M A_0 \) must be of order of \( T_{TJ} \sim M_{H_0}/25 \). As the mass of \( H_0 \) goes above \( W \), \( Z \) or \( h \) threshold, \( H_0 \) annihilation into \( WW \), \( ZZ \) and \( hh \) become increasingly efficient, an effect which strongly suppresses the \( H_0 \) relic density. The region 4 of figure 1, corresponding to \( M_{H_0} \in [40, 80] \) GeV, appears to be the only region consistent with WMAP data.

For the high mass regime, we can derive the general trends from figure 1, the plot on the right. No new annihilation channel opens if \( M_{H_0} \) is heavier than the Higgs or the gauge bosons. There are then essentially two kinds of processes which control both the abundance: the annihilation into two gauge bosons, dominant if \( \mu_2 < M_{H_0} \), and the annihilation into two Higgs, which dominates if \( \mu_2 > M_{H_0} \). Coannihilation plays little role.

The abundance of dark matter is suppressed over most of the area of the plot because of large quartic coupling effects on the cross-sections. Let us emphasize that in this regime, it is

---

\( H_0 H^+ \) coannihilation is suppressed for our choice of \( \Delta M H_c \).
Figure 2: Left: Relic density and Right: scattering cross-section intervening in direct detection searches, all as a function of the mass of dark matter and comparison with the MSSM. For the direct detection plot, the light colors correspond to $0.01 < \Omega_{DM} h^2 < 0.3$, while the dark colors correspond to $0.094 < \Omega_{DM} h^2 < 0.129$.

possible to reach agreement with WMAP data, but only at the price of some fine tuning. We need to keep the mass splittings between the components of $H_2$ relatively small. First because large mass splittings correspond to large couplings and second because the different contributions to the annihilation cross-section must be suppressed at the same location, around $\lambda_L - 0$ (i.e. $M_{H_0} \simeq \mu \simeq M_{A_0} \simeq M_{H_+}$ in this case). As it can be seen in figure 1, the plot on the right, the area consistent with WMAP corresponds to the narrow region around the diagonal with $M_{H_0} \gtrsim 800$ GeV for $\Delta M_{A_0} = 5$ GeV and $\Delta M_{H_+} = 10$ GeV. Notice that this behavior is limited by the unitarity bound on the total annihilation cross-section\(^\text{10}\) which constrains the mass of the dark particle to be $M_{H_0} \lesssim 120$ TeV.

In figure 2, the plot on the left, we show a scatter plot of $\Omega_{DM} h^2$ as a function of the mass of the dark matter candidate $M_{DM}$ for a fair sample of IDM (scanning on several Inert scalar mass splittings) and, for the sake of comparison, for the MSSM. We clearly see the two regimes (low mass and high mass) of the IDM that may give rise to a relevant relic density (i.e. near WMAP). The MSSM models have a more continuous behavior, with $\mathcal{O}(100 \text{ GeV})$ dark matter masses. As a conclusion, the IDM provides dark matter candidates with masses as small as 40 GeV and as large as 600 GeV in contrast with the MSSM more concentrated around $\sim 100$ GeV.

4 Direct detection

Direct detection searches look for signals of dark matter in low background detectors trying to measure the energy deposited by the scattering of a dark matter particle with a nucleus of the detector. Assuming that the main interaction contributing to the $H_0$-quarks interaction is the spin independent $H_0 q^h \rightarrow H_0 q^h$ interaction\(^\text{1}\) it can be shown\(^\text{2}\) that the $H_0$ elastic scattering cross section off a proton scales like

$$\sigma_{H_0-p} \propto \lambda_L^2/(M_{H_0} M_h^2)^2.$$  \(\text{(3)}\)

In figure 2, the plot on the right, we show a scatter plot of $\log_{10} \sigma_{DM-p}$ as a function of $M_{DM}$.

\(^\text{1}\)The experiments have reached such a level of sensitivity that the $Z$ exchange contribution $H_0 q^h \rightarrow A_0 q^h$ is excluded by the current experimental limits\(^\text{4}\). Consequently, to forbid $Z$ exchange by kinematics, the mass of the $A_0$ particle must be larger than the mass of $H_0$ by a few 100 keV.
for the IDMs considered in the abundance plot of figure 2 that account for $0.01 < \Omega_{DM} h^2 < 0.3$. We see that the low mass regime candidates could be detected by future experiments such as EDELWEISS II or by the ton sized experiments such as ZEPLIN. For the higher mass regime however, there is no hope for future detection in low background detector. Indeed, the WMAP requirement for dark matter relic density constrains the $\lambda_L$ couplings to be vanishing while the same couplings drive the amplitude of the matter-$H_0$ scattering cross-section.

5 Indirect detection

The measurement of secondary particles coming from dark matter annihilation in the halo of the Galaxy is another promising way of deciphering the nature of dark matter. Let us emphasize that this possibility depends however not only on the properties of the dark matter particle, through its annihilation cross-sections, but also on the astrophysical assumptions made concerning the distribution of dark matter in the halo that supposedly surrounds our Galaxy.

In Figure 3, we show the log of the produced gamma-ray flux from dark matter annihilation at the Galactic Center $\Phi_\gamma$ as a function of $M_{DM}$ for the same sample of models than for direct detection. We computed the gamma-ray flux for the plot on the left assuming an isothermal dark matter density profile while for the plot on the right we assumed a Navarro, Frank and White (NFW) profile. The main difference between these two profiles is the slope of the dark matter density as a function of the galactic radius in the central part of the Galaxy. The isothermal profile is flat while the NFW profile is more cuspy (i.e. steeper). We see that for steeper profile the gamma-ray flux is larger.

The particle physics dependence of $\Phi_\gamma$ also clearly show up in figure 3. Indeed we see that $\Phi_\gamma$ behaves differently in the low and the high mass regime of the IDM given that the processes contributing to the annihilation cross-section are different. Moreover, notice that the IDM dark matter candidates have typically higher detection rates than the neutralino in SUSY models, especially at high mass. Let us stress that the figures for indirect detection were obtained taking into account annihilation processes at three level only (see 4, for a recent study of the IDM...
including processes at one-loop). It can be inferred from figure 3 that the IDM can give the right relic abundance in a range of parameters which will be probed by GLAST for NFW dark matter profiles. GLAST will however have no chance to observe the gamma-ray flux produced by annihilating $H_0$ at the Galactic Center for flatter profiles such as the isothermal one.

6 Conclusion

We carried out a rather detailed analysis of the IDM as a dark matter model assuming the standard freeze-out mechanism. We recovered the results of Barbieri et al. and Cirelli et al. which a priori did not seem to match. This is because the IDM provides dark matter candidates in two rather separate mass ranges, one between 40 and 80 GeV, the other one over 400 GeV. The physics driving the existence of dark matter in these regions of the parameter space is quite different.

We have also investigated the prospects for direct and indirect detection searches. Concerning direct detection searches, the low mass regime candidates should be detected with the future ton sized experiments while the high mass regime will stay out of reach. For indirect detection searches we looked at the gamma-ray flux generated at the Galactic Center by dark matter annihilation. Whatever the dark matter density profile assumed, we have come to the conclusion that the Inert scalars have typically higher detection rates than the neutralino in SUSY models, especially at high mass. Moreover, the IDM could be probed by the future GLAST experiment.

Acknowledgments

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COSMOLOGICAL CONSTRAINTS ON GAUGINO MEDIATION

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With gaugino-mediated supersymmetry breaking, the gravitino, a neutralino or a scalar lepton can be the lightest or next-to-lightest superparticle. We discuss cosmological constraints on the different scenarios. While neutralino dark matter and gravitino DM with a sneutrino NLSP are consistent for a wide range of parameters, gravitino DM with a stau NLSP is strongly constrained. Gravitino DM with a neutralino NLSP is excluded.

1 Introduction

Gaugino-mediated supersymmetry breaking\(^2^,\,^3\) employs a setup with extra space-time dimensions to generate viable soft masses for the superparticles. SUSY is broken by the \(F\)-term vacuum expectation value of a field which is localised on a brane. The MSSM squarks and sleptons live on another brane located at a different position in the extra dimensions. Therefore, they cannot couple to the SUSY-breaking field and do not obtain soft masses at the tree level. Gauginos and Higgses are assumed to be bulk fields, so that they receive soft masses via direct interactions. Thus, assuming that the gauge couplings unify and that the compactification scale of the extra dimensions is close to the unification scale, one obtains at this energy

\begin{equation}
g_1 = g_2 = g_3 = g \approx \frac{1}{\sqrt{2}},
\end{equation}

\begin{equation}
M_1 = M_2 = M_3 = m_{1/2},
\end{equation}

\begin{equation}
m_{\tilde{\nu}_L}^2 = m_{\tilde{\nu}_R}^2 = 0 \quad \text{for all squarks and sleptons } \tilde{\phi},
\end{equation}

\begin{equation}
A_{\tilde{\phi}_L} = 0 \quad \text{for all squarks and sleptons } \tilde{\phi},
\end{equation}

\begin{equation}
\mu, B\mu, m_{h_i}^2 \neq 0 \quad (i = 1, 2),
\end{equation}

\(^{\text{a}}\)Talk presented at the XLInd Rencontres de Moriond, March 10-17, 2007, La Thuile, Italy. Based on work done in collaboration with Wilfried Buchmüller, Laura Covi and Kai Schmidt-Hoberg.\(^1\).
where the GUT charge normalisation is used for $q_1$ and where $m_{h_1}^2$ and $m_{h_2}^2$ are the soft masses of the Higgs fields coupling to the down-type and the up-type quarks, respectively. For simplicity, we restrict ourselves to the case $m_{h_2}^2 \geq 0$. A discussion that includes negative values can be found in $^4$.

The renormalisation group running from the compactification scale to low energies produces non-zero soft masses for all superparticles. We employed SOFTSUSY $^5$ to calculate the mass spectrum. As benchmark values for our discussion, we chose $m_{1/2} = 500$ GeV, $\tan \beta = 10$ or 20 and $\text{sign}(\mu) = +1$. The values of $\mu$ and $B\mu$ are determined by the conditions for electroweak symmetry breaking. Due to the large effects of the strong interaction, the squark masses experience the fastest running and end up around a TeV. Regarding the mass ordering of the sleptons and neutralinos, the most important free parameters are the soft Higgs masses. Depending on their values, the lightest MSSM superparticle can be a neutralino, a stau or a sneutrino $^{2,3,6}$. Besides one of these particles, the gravitino can be the lightest superparticle (LSP). In gaugino mediation its mass is bounded from below $^7$. We use $m_{3/2} \geq 10$ GeV, relaxing the limit from naive dimensional analysis by a factor 5 in order to obtain a conservative lower bound.

The constraints on the parameter space of gaugino mediation that arise from collider experiments were described in $^{8,3,4}$. In the following, we will present the impact of constraints from cosmology, summarising and updating the study $^1$.

2 Cosmological Constraints

As long as $R$ parity is conserved, the LSP is stable and its energy density has to be consistent with the observed cold dark matter density. We use the $3\sigma$ range given in $^{10}$,

$$0.106 < \Omega_{\text{DM}} h^2 < 0.123 \ . \ (6)$$

If the gravitino is the LSP, the next-to-lightest superparticle (NLSP) decays during or after big bang nucleosynthesis (BBN). Analogously, the gravitino decays late if it is not the LSP. The very energetic decay products of such long-lived particles can alter the primordial light element abundances compared to the standard BBN scenario. This leads to constraints on the released electromagnetic and hadronic energy $^{11,12,13,14}$, which can be quantified approximately by upper bounds on the product $\epsilon_{\text{em,had}} Y_{\text{NLSP}}$. Here $\epsilon_{\text{em,had}}$ is the average electromagnetic or hadronic energy emitted in a single NLSP decay, and the abundance $Y_{\text{NLSP}}$ is given by the NLSP number density prior to decay divided by the total entropy density. The abundance was determined numerically using micrOMEGAs $^{15,16}$ and assuming that the NLSP freezes out with its thermal relic density. We used the bounds on the energy release compiled in $^{17}$, which were computed in the earlier studies $^{11,12}$. These bounds were derived from the observed abundances of deuterium and $^4$He. Following $^{17}$, we considered two sets of constraints in order to take into account the uncertainties in the measurements. Points in parameter space violating the “conservative” limits are excluded, while points violating only the “severe” limits should be regarded as disfavoured. In addition, we applied the constraint derived from the ratio of the observed $^3$He and D abundances in $^{13}$.

Decays of long-lived particles in the early universe can also cause distortions of the cosmic microwave background. However, the corresponding constraints on late electromagnetic energy release are less constraining than the BBN bounds in our case $^{17}$.

Recently it was pointed out that negatively charged stau NLSPs can form bound states with light nuclei, which leads to a drastic change of the reaction rates relevant for BBN $^{18}$. In $^{18}$ and in the subsequent studies $^{19,20,21}$ it was found that this effect causes an unacceptable overproduction of $^6\text{Li}$ unless the stau lifetime is less than $10^3 - 10^4$ s.
3 Neutralino Dark Matter

Fig. 1 shows the part of the allowed parameter space for the soft Higgs masses where a neutralino is the lightest MSSM superparticle for $\tan \beta = 20$. Its mass varies between 83 GeV and 204 GeV. If this neutralino is also lighter than the gravitino, it can form the dark matter. In the black region of the figure, the thermal neutralino relic density lies within the experimentally allowed range (6). The bino contributes at least 80% to the lightest neutralino. The remaining components are chiefly Higgsinos. Towards the left end of the region, the lightest neutralino is a pure bino. In the magenta (dark-gray) region in Fig. 1, the thermal neutralino density is smaller than the lower bound in Eq. (6) and hence only constitutes a part of the dark matter density. This is not ruled out, since the dark matter could be made up of several components or of both thermally and non-thermally produced neutralinos. The lightest neutralino is an almost pure Higgsino at the right edge of the parameter space. Finally, the upper limit on the dark matter density excludes the white region in Fig. 1. For $\tan \beta = 10$, the results are similar.

The cross-section for the direct detection of neutralino dark matter is suppressed for a pure bino. Therefore, in the largest part of the allowed parameter space direct dark matter detection will be very hard. The cross-section is larger towards the right end of the region, where the lightest neutralino has a significant Higgsino component. Although the cross-section is still at least one order of magnitude below the present bounds there\textsuperscript{22,4}, it could be accessible to the next generation of dark matter experiments.

In the neutralino LSP scenario, BBN leads to an upper bound on the density of gravitinos, since their decays affect the light element abundances as explained above. This bound can be translated into an upper limit on the reheating temperature after inflation\textsuperscript{11,12}. The latter is a free parameter in our discussion and therefore it can simply be assumed to be sufficiently small. The other superparticles decay into the LSP before the start of BBN and thus do not cause problems either, unless LSP and NLSP are nearly degenerate. Consequently, neutralino dark matter is a viable scenario in gaugino mediation.
Gravitino Dark Matter

If the gravitino is the LSP, it is a viable dark matter candidate as well. The thermally produced gravitino density falls into the range $6$ for a reheating temperature between about $10^9$ GeV and $10^8$ GeV, if $10\text{ GeV} \lesssim m_{3/2} \lesssim 100 \text{ GeV}$. Non-thermal production from NLSP decays is negligible in the scenario under consideration.

A neutralino NLSP together with a gravitino LSP heavier than a GeV is excluded by the hadronic BBN constraints for $m_Z > m_{3/2} + m_Z$, since in this case two-body decays can produce $Z$ bosons, which have a large hadronic branching ratio $^{23,24}$. In the parameter space region where the neutralino is so light that the decay into real $Z$ bosons is kinematically forbidden, we find abundances that are small enough to satisfy the hadronic bounds. However, the electromagnetic constraints turn out to be violated, so that gravitino dark matter with a neutralino NLSP is not allowed in gaugino mediation.

If a scalar tau is the NLSP, the hadronic energy release from stau decays is always in agreement with the BBN bounds. However, the bounds on the electromagnetic energy release $e_{\text{en}} Y_{\tau}$ are significantly more constraining. They are shown in Fig. 2 as functions of the stau lifetime $\tau_{\tau}$. The curves exclude (or disfavour in the case of the severe bound from D) the area above them. The curve derived from the $^3\text{He}$ abundance is only shown in the region where it is more constraining than the conservative deuterium bound. The limit arising from the $^4\text{He}$ abundance is not shown, since it is not relevant in the stau NLSP region.

We also plot points from the $\tilde{\tau}$ NLSP region for $\tan \beta = 10$ in the $e_{\text{en}} Y_{\tau} - \tau_{\tau}$ plane, assuming that half the energy of the tau produced in the dominant two-body stau decay contributes to the electromagnetic energy release. The remaining energy ends up in neutrinos and is not relevant. The red (dark-gray) points are the results for a gravitino mass of $50 \text{ GeV}$. We find that the conservative $D$ constraints can be satisfied. In order to satisfy the $^3\text{He}$ and the severe $D$ bounds, the NLSP lifetime has to be shorter. This is the case for smaller gravitino masses. A rough estimate of the upper limit from the $^3\text{He}$ bound is $m_{3/2} \lesssim 30 \text{ GeV}$. For $m_{3/2} = 10 \text{ GeV}$, we see from the green (light-gray) points in Fig. 2 that most of the stau NLSP region is allowed by all constraints derived from the deuterium and helium abundances.

However, it is also obvious that the stau lifetime violates the upper bound $\tau_{\tau} \lesssim 10^3 - 10^4 \text{ s}$ that arises when bound state effects are included in the BBN calculations $^{18,19,20,21}$. We find $86 \text{ GeV} \leq m_{\tilde{\tau}} \leq 203 \text{ GeV}$ in the $\tau$ NLSP region, which corresponds to $\tau_{\tau} \geq 1.8 \cdot 10^5 \text{ s}$. Thus, we have to conclude that BBN disfavours gaugino mediation with a gravitino LSP and a stau NLSP for our benchmark value $m_{1/2} = 500 \text{ GeV}$. In order to decrease the stau lifetime sufficiently, one has to increase the superparticle masses roughly by a factor 3. The constraint could also be avoided if there was sizable entropy production between the decoupling of the staus from the thermal bath and the start of BBN $^{25,20}$. The last NLSP candidate in gaugino mediation is the sneutrino. In the corresponding parameter space region, we find sneutrino masses and lifetimes which lie roughly in the same range as those of the staus. The relic abundance is also similar, $1.3 \cdot 10^{-13} \leq Y_{\tilde{\nu}} \leq 4.6 \cdot 10^{-13}$. The BBN bounds on a sneutrino NLSP are rather weak, since the dominant two-body decay produces gravitinos and neutrinos. These neutrinos can release electromagnetic energy by annihilating with background (anti-)neutrinos and producing charged particles. Besides, hadronic constraints arise from the four-body decay into a neutrino, a gravitino and a quark-antiquark pair. They turn out to be more stringent but allow thermally produced sneutrinos with masses up to around $300 \text{ GeV}$ according to $^{14}$. This result is based on an approximation for the sneutrino abundance. In our case, the abundance turns out to be larger due to the importance of co-annihilations. Hence, a part of the sneutrino region may be in conflict with BBN $^{26}$. However, one has to bear in mind that $^{14}$ assumed a much smaller experimental uncertainty for the deuterium abundance than the studies $^{13,17}$, on which our analysis of the stau NLSP scenario is based.
5 Conclusions

We have discussed dark matter candidates in theories with gaugino-mediated supersymmetry breaking. The parameters relevant for the superparticle mass spectrum are the universal gaugino mass, the soft Higgs masses, $\tan \beta$ and the sign of $\mu$. At different points in parameter space, the gravitino, a neutralino or a slepton can be the lightest or next-to-lightest superparticle.

We have investigated constraints from the observed dark matter density and big bang nucleosynthesis on the different scenarios. A neutralino LSP as the dominant component of dark matter is a viable possibility. Gravitino dark matter with a neutralino NLSP is excluded. A stau NLSP can be consistent with the constraints arising from the effects of stau decays on the primordial light element abundances. However, taking into account bound state effects on the lithium abundance disfavour this scenario. Gravitino dark matter with a sneutrino NLSP can also be realised and is nearly unaffected by all constraints.

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Natural SUSY Dark Matter: A Window on the GUT Scale

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One of the key motivations for supersymmetry is that it provides a natural candidate for dark matter. For a long time the density of this candidate particle fell within cosmological bounds across much of the SUSY parameter space. However with the precision results of WMAP, it has become apparent that the majority of the SUSY parameter space no longer fits the observed relic density. This has given rise to claims that supersymmetry no longer provides a natural explanation of dark matter. We address this claim by quantifying the degree of fine-tuning required for the different dark matter regions. We find that the dark matter regions vary widely in the degree of tuning required. This degree of tuning can then be used to provide valuable insights into the structure of SUSY breaking at the GUT scale.

1 Introduction

Supersymmetry at the TeV scale is one of the most compelling candidates for physics beyond the Standard Model (SM). A primary motivation for supersymmetry is that it removes the need to fine-tune the bare Higgs mass. It also naturally provides a candidate for cold dark matter. If we are to avoid fast proton decay, we must introduce a symmetry that constrains the interactions of particles with their supersymmetric partners. The most common form of this symmetry is $R$-parity. This forbids the decay of a single superpartner into purely SM matter. One result of this is that the lightest superparticle (LSP) is absolutely stable. If SM matter and superparticles were in thermal equilibrium in the early universe, the cooling universe would leave behind a relic density of superparticles.

This has given rise to many claims that SUSY naturally accounts for dark matter. However, having a candidate particle is one thing whereas naturally accounting for the observed relic density is quite another. In fact as WMAP has improved the constraints on the relic density the regions of the SUSY parameter space that fit the observed relic density have begun to look
very slender. This has led to recent claims that low energy supersymmetry requires significant fine-tuning to fit the data that others claim it accounts for ‘naturally’.

One could ignore such a war of words over what is or is not natural. However SUSY derives a significant portion of its motivation from questions of tuning and naturalness. Therefore this question deserves to be taken seriously. Here we present a quantitative study of the fine-tuning required to access the different dark matter regions of the Minimal Supersymmetric Standard Model (MSSM). We discuss the implications of these different degrees of tuning for the MSSM. Finally, we highlight how considerations of tuning allow us to compare GUT scale models of SUSY breaking from LHC data.

2 Fine-tuning and Dark Matter

To quantitatively study fine-tuning we need a measure. The fine-tuning required for electroweak symmetry breaking has a long history of quantitative study. We follow Ellis and Olive in using an analogous measure to study the fine-tuning of dark matter:

\[ \Delta^\Omega_a = \left| \frac{\partial \ln (\Omega_{CDM} h^2)}{\partial \ln (a)} \right| \]

(1)

where \( \{a\} \) are the free parameters of the theory. This provides a measure of the sensitivity of the dark matter relic density to the inputs. For example, if \( \Delta^\Omega_a = 10 \), a 1% variation in \( a \) would result in a 10% variation in \( \Omega_{CDM} h^2 \).

3 The Constrained Minimal Supersymmetric Standard Model (CMSSM)

The MSSM is notorious for having over 100 free parameters. However, many of these parameters are already constrained by experiment to be zero. Furthermore, the parameters are only free if we leave the mechanism of SUSY breaking entirely unspecified. In a realistic theory, we would expect all the MSSM parameters to be set in terms of a smaller set of more fundamental parameters. In the absence of a specific theory of SUSY breaking we can still make progress. The most frequently studied SUSY model is the CMSSM with the parameters:

\[ a_{CMSSM} \in \{ m_0, \ m_{1/2}, \ A_0, \ \tan \beta \ and \ sign(\mu) \} \]

(2)

Here \( m_0 \) is the common soft mass of all the scalar particles, \( m_{1/2} \) is the common mass for all the gauginos, \( A_0 \) sets the third family trilinear couplings, \( \tan \beta \) is the ratio of the two Higgs vevs and \( \mu \) is a bilinear Higgs mass term.

In Fig. 23 we show the \((m_{1/2}, m_0)\) plane of the CMSSM parameter space for \( \tan \beta = 10 \), \( A_0 = 0 \) and \( \mu \) positive. Low \( m_0 \) is ruled out (light green) as it results in a \( \tilde{\tau} \) LSP. This would result in a charged relic which would have been observed in searches for anomalously heavy nuclei. Low \( m_{1/2} \) is ruled out (light blue) as this results in a light Higgs \( m_h < 111 \text{ GeV} \). In the remaining parameter space we plot the SUSY contribution to \( (g-2) \) of the muon, \( \delta a_\mu \). We take the current observation of a deviation from the Standard Model seriously and plot the region in which we agree with the measurement at \( 2\sigma \) (long dashed green lines) and \( 1\sigma \) (short dashed green lines). It is clear that to explain the observed value of \( \delta a_\mu \) we require light soft SUSY masses and thus light superpartners. Finally we plot the band that fits the observed relic density

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of dark matter within 2$\sigma$. For every point that lies in this band we calculate the fine-tuning and plot the point in a colour that corresponds to the log-scale on the right of the plot.

The only dark matter region in Fig. ?? lies alongside the region in which the LSP is the $\tilde{\tau}$. Above this region the LSP is the bino (the partner to the $U(1)$ gauge boson of the Standard Model). Bino LSPs annihilate very weakly and normally give $\Omega_{\text{CDM}} h^2 \gg \Omega_{\text{CDM}}^{\text{WMAP}} h^2$. Thus, above the dark matter strip we have too much dark matter and WMAP rules out the CMSSM. The remaining parameter space is very slender. In the remaining parameter space, the $\tilde{\tau}$ and the lightest neutralino, $\tilde{\chi}_1^0$, are very close in mass. This results in a large number density of both particles in the early universe. The resulting coannihilation of staus and neutralinos greatly enhances the annihilation rate of SUSY matter, greatly lowering the resulting relic density. This is very sensitive to the mass difference between the stau and the lightest neutralino. Thus we would expect a coannihilation region to be fine-tuned. However note that the band is green at low $m_{1/2}$ and red at large $m_{1/2}$. This corresponds to a tuning of $3 - 10$. This is a surprisingly small degree of fine-tuning. The reason is that, for low $m_0$ and low tan$\beta$, the mass of the stau and the lightest neutralino are both primarily dependent on $m_{1/2}$. Therefore the masses of the coannihilating particles are coupled and the majority of the fine-tuning is removed.

This is a special case. In Fig. ?? we show the $(m_{1/2}, m_0)$ plane of the CMSSM with tan$\beta = 50$, $A_0 = 0$ and $\mu$ positive. We have also extended the range of $m_0$ and $m_{1/2}$. Many of the bulk features remain the same. Low $m_0$ is still ruled out by a $\tilde{\tau}$ LSP. Low $m_0$ and $m_{1/2}$ is ruled out by a light Higgs. A new bound rules out large $m_0$ and low $m_{1/2}$ (light red). Here the parameter space gives $\mu^2 < 0$ which is unphysical and corresponds to a failure of radiative electroweak symmetry breaking (REWSB).

The dark matter phenomenology is noticeably more complex. As before, the LSP is bino across the majority of the parameter space and thus mostly gives $\Omega_{\text{CDM}} h^2 \gg \Omega_{\text{CDM}}^{\text{WMAP}} h^2$. The exceptions to this are marked by the thin WMAP strips. Once again we have a coannihilation strip lying along the side of the $\tilde{\tau}$ LSP region. In contrast to Fig. ??, this band is plotted in purple, corresponding to $\Delta\Omega \approx 50$.

This band is broken by two bands that go up in both $m_0$ and $m_{1/2}$. These lie on either side of the line along which $2m_{\chi_1^0} = m_A$ and the neutralinos annihilate via an s-channel pseudoscalar.
Higgs boson. As could be expected, such a process enormously enhances the annihilation of dark matter. We only fit the dark matter relic density with just enough resonant annihilation. This sounds like fine-tuning and indeed the lines are mostly plotted in grey indicating $\Delta \Omega > 100$.

Finally there is a dark matter band that runs alongside the region in which $\mu^2 < 0$. Along the edge of this region we have low $\mu$ and the higgsino fraction of the lightest neutralino increases. As bino dark matter gives $\Omega_{CDM} h^2 \gg \Omega_{CDM} WMAP h^2$, and higgsino dark matter generally gives $\Omega_{CDM} h^2 \ll \Omega_{CDM} WMAP h^2$, it is not surprising that somewhere in between we manage to fit the observed dark matter relic density. However, the requirement that the composition of the LSP include just enough higgsino is an indication of fine-tuning and indeed the line is plotted in purple and red indicating a tuning $\Delta \Omega \approx 30 - 60$.

4 Breaking the Constraints

We have shown the typical tunings of dark matter in the CMSSM, and highlighted the problem that the bino LSP results in $\Omega_{CDM} h^2 \gg \Omega_{CDM} WMAP h^2$, ruling out the majority of the CMSSM parameter space. However, there are few compelling theoretical reasons to remain within the confines of the CMSSM. Indeed there are many good reasons to relax a number of the constraints. In previous work\(^7\) we study the implications for fine-tuning of relaxing the constraint of universal scalar masses and universal gaugino masses. In Fig. ?? we consider a model in which we allow the gaugino masses to vary independently of one another. Such a model has the parameters:

$$a_{CMSSM+M} \in \{m_0, M_1, M_3, M_3, A_0, \tan \beta \text{ and } \text{sign} (\mu)\}$$

(3)

where $M_{1,2,3}$ set the GUT scale soft SUSY breaking mass of the superpartners to the $U(1)$, $SU(2)$ and $SU(3)$ gauge bosons respectively. Such a break from gaugino mass universality can arise naturally in string models\(^7\) and GUT models\(^7\).

In Fig. ??, we show the $(M_1, m_0)$ plane of a model with non-universal gaugino masses where we have fixed $M_{2,3} = 350$ GeV, $A_0 = 0$ and $\tan \beta = 10$ with $\mu$ positive. As before there is a region at low $m_0$ that is ruled out by a $\tilde{\tau}$ LSP. There is a region ruled out at low $m_{1/2}$ due to light neutralinos and a region ruled out at light $m_0$ and $m_{1/2}$ due to light sleptons.
Figure 3: The \((\tilde{m}_1, m_0)\) plane of the CMSSM \(+M_1\) parameter space with \(M_{2,3}=350\) GeV, \(\tan\beta=10\), \(A_0=0\) and \(\mu\) positive.

The most notable feature is the explosion in the complexity of the dark matter regions. Now, rather than the three dark matter regions of the CMSSM, we have five distinct dark matter regions. Firstly, there is the familiar bad along the edge of the \(\tilde{\tau}\) LSP region due to \(\tilde{\tau}-\chi_1^0\) coannihilation. As before, this exhibits low fine-tuning. This band is interrupted at \(M_1 \approx 570\) GeV. At larger \(M_1\) the neutralino is wino rather than bino. A wino LSP generally gives \(\Omega_{CDM} h^2 \ll \Omega_{CDM}^{WMAP} h^2\). For \(M_1 < 580\) GeV the neutralino is bino so \(\Omega_{CDM} h^2 \gg \Omega_{CDM}^{WMAP} h^2\). Around \(M_1 \approx 570\) GeV the neutralino has just enough bino and wino to fit the observed dark matter density. As in the mixed bino/higgsino case, this requires a delicate balance and thus the region exhibits a tuning \(\Delta^\Omega \approx 30\).

At low \(M_1\) there are two distinct peaks at \(M_1 = 110\) GeV and \(M_1 = 130\) GeV. These correspond to neutralino annihilation via an s-channel \(Z\) or light Higgs boson respectively.

Finally, there is a wide band that fits the observed dark matter relic density at low \(m_0\) and low \(m_{1/2}\). It lies alongside the region that is ruled out by LEP searches for light sleptons. In this band, the sleptons are light enough to enhance the decay of neutralinos via t-channel slepton exchange to the point where we suppress the dark matter relic density enough to fit the observed data. This decay process is remarkably insensitive to the precise value of the soft masses. This translates to \(\Delta^\Omega < 1\), corresponding to no fine-tuning.

Fig. ?? presents one example of the dark matter phenomenology of the wider MSSM. By relaxing the constraints of the CMSSM we can find the typical tunings of the different dark matter regions that exist within the MSSM. We list these in Table ?? . Thus we conclude that each region has a typical tuning, and that there remain regions of the MSSM that require no tuning to accommodate the observed relic density.

5 Conclusions: interpreting fine-tuning

We must be careful in our interpretation of these results. Just because an MSSM dark matter region exhibits significant fine-tuning, does not mean that such a region will not be found at a future collider. These tunings have been calculated with respect to the MSSM, which is an
Table 1: The typical tunings for dark matter regions within the MSSM.

<table>
<thead>
<tr>
<th>Region</th>
<th>Typical $\Delta^{14}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mixed bino/wino</td>
<td>$\sim 30$</td>
</tr>
<tr>
<td>Mixed bino/higgsino</td>
<td>$30 - 60$</td>
</tr>
<tr>
<td>Mixed bino/wino/higgsino</td>
<td>$4 - 60$</td>
</tr>
<tr>
<td>Bulk region (t-channel $f$ exchange)</td>
<td>$&lt; 1$</td>
</tr>
<tr>
<td>slepton coannihilation (low $M_1$, $m_0$)</td>
<td>$3 - 15$</td>
</tr>
<tr>
<td>slepton coannihilation (large $M_1$, $m_0$, $\tan \beta$)</td>
<td>$\sim 50$</td>
</tr>
<tr>
<td>$Z$-resonant annihilation</td>
<td>$\sim 10$</td>
</tr>
<tr>
<td>$h^0$-resonant annihilation</td>
<td>$10 - 1000$</td>
</tr>
<tr>
<td>$A^0$-resonant annihilation</td>
<td>$80 - 300$</td>
</tr>
</tbody>
</table>

**effective theory.** The MSSM does not specify the mechanism of SUSY breaking, instead we parameterise our ignorance with soft SUSY breaking masses. We expect that these masses should be set by a deeper theory. Thus a region that is tuned in the MSSM may have a very different tuning within a specific model of SUSY breaking.

This variation of tuning between models provides us with a useful tool. If the LHC finds signals for new physics in the form of new particles and large quantities of missing energy, many will interpret this as a SUSY mass spectrum. There will be many different high energy models that fit the data, and probably many models that will also fit the observed dark matter relic density. However this raw mass spectrum will do little to tell us what relations must obtain between high energy parameters.

If we analyse the sensitivity of the dark matter relic density in such a scheme we test precisely this dependence between high energy parameters. For example, it is only because the $\tilde{\tau}$ mass and the $\chi^0_1$ mass in the CMSSM coannihilation region are both dominantly set by $m_{1/2}$ that such a region has low tuning. Therefore we would have to favour such an explanation of an observed coannihilation region than a model that set both masses independently.

After we identify the relations that mitigate the fine-tuning, we can go on to make further predictions. Thus questions of fine-tuning can help to narrow down the candidate explanations for a given experimental signal. Having done this, novel predictions can be made on the basis of hypothesised GUT scale relations between the soft masses, and these can be tested in future experiments. Thus fine-tuning and naturalness should allow us to analyse and compare GUT scale physics using LHC energy data.

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**References**

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In this letter in the framework of a simple see-saw scenario with three quasi degenerate Majorana neutrinos we propose that one of these neutrinos can be very weakly coupled, yet there is a mechanism of the generation of the abundance of such "dark" neutrino in the early universe. The mechanism of production is due to oscillations between "dark" Majorana neutrino and one of the Majorana's which has relatively large Yukawa couplings ("bright" Majorana neutrino). The transition of "bright" Majorana into a "dark" one is nonadiabatic. We point out on the similarity with the Landau-Zener transition regime. In our model one can explain the observed dark matter density, present matter-antimatter asymmetry and active neutrino data all at the same time.

1 Introduction

The observed matter-antimatter asymmetry and the presence of yet unknown form of matter in the universe known as Dark Matter (DM) both require the extension of the Standard Model which does not contain neither a DM candidate nor a mechanism to produce large enough baryon asymmetry in order to match the observed value. On the other hand, the observed neutrino mass scales and mixings can not be incorporated within the Standard Model (SM) as well. The sew-saw mechanism\(^2\) provide probably the most natural and simple explanation why neutrinos are massive, yet so light compared to all other Standard Model particles. At the same time, neutrino mixing data represent a positive test for leptogenesis\(^3\), an attractive explanation of the the observed matter-antimatter asymmetry of the Universe as a direct consequence of the see-saw mechanism. Can DM be accommodated within the see-saw model as well? It is known, that it is enough to have just two Majorana's to account for both leptogenesis and active neutrino data (see, e.g. \(^4\)). Assuming that the number of Majorana neutrinos is the same as the number of lepton flavors in Standard Model one can ask the question if one of the Majorana neutrino can
be made almost decoupled and, thus, stable at present but being produced efficiently enough in the early universe. However, if a particle is almost stable at present there is typically quite hard to produce it during earlier epoch. One clear possibility is the DM production from the inflaton decay along with the rest of particles. There is no argument why not to assume, for example, that inflaton couplings to the DM particle and to the SM matter are not of the same order. In that case DM particle can be produced as efficiently as ordinary matter. However, such mechanism, due to the present lack of the detailed picture of inflation and most importantly the exit from it, is strongly model dependent. Besides, nothing would require one of the Majorana neutrinos to be a DM particle. Any other particle decoupled from the SM fields would be as plausible. It is appealing, therefore, to look for another more constrained mechanism of the production of the "dark" Majorana neutrino. The perturbative production out of the thermal bath can be a solution. However, typically this would require the reheating temperature after inflation to be quite high.

In this letter we propose that "dark" Majorana neutrino is produced via oscillation mechanism which is similar to the Landau-Zener type transition. The production starts and terminates near some critical temperature $T_c$ due to the oscillations between "dark" Majorana neutrino and one of the strongly coupled (or "bright") Majorana neutrino. The critical temperature can be as low as the Majorana neutrino mass scale which, in turn, can be as low as $\mathcal{O}(\text{TeV})$. Contrary to the perturbative production mechanism in our model high reheating temperature is not required. In order for oscillation mechanism to be successful both neutrinos have to be quasi degenerate in mass. We will discuss two scenarios. The one, called "minimal", has only the see-saw type couplings being added to the SM Lagrangian. In this case the production rate turns out to be not efficient enough. In the "next-to-minimal" scenario we will introduce a simple dimension five operator which couples Higgs field to Majorana neutrinos and is suppressed by the high energy physics scale $\Lambda$. Such operator could result from GUT extensions of the SM, supergravity etc. This operator is nearly irrelevant for any physics at present. However, at relatively high temperatures ($T > T_{\text{EW}}$) it will turn out to be an effective source for the "dark" Majorana production.

2 The Model.

The see-saw extension of the Standard Model is described by the Lagrangian

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + i\bar{N}_k \partial \bar{N}_k - F^{ik} \bar{L}_i N_k \phi^c - \frac{M}{2} \bar{N}_i^c N_i + \text{h.c.}$$

For simplicity we will assume that only one of the "bright" Majorana neutrino serves as a source for the production of the "dark" Majorana neutrino. We will employ the density matrix formalism\textsuperscript{10,11}. The evolution of the number densities of the Majorana and of the active neutrinos is described by the density matrix which obeys quantum kinetic equation:

$$i\dot{\rho}_k = [H_k, \rho_k] - i\frac{\Gamma_k}{2} (\Gamma^+_k + \Gamma_k^-) \rho_k + i\Gamma^+_k$$

\textsuperscript{a}In the case of quasi degenerate neutrinos the Majorana mass scale can be made quite low as the consequence of resonant enhancement of the kinematic part of the expression for the CP-asymmetry, which compensates for the smallness of the Yukawa couplings\textsuperscript{a}.

\textsuperscript{a}A simultaneous picture of baryogenesis and DM has been already proposed\textsuperscript{5,7}. Here the matter-antimatter asymmetry is generated through RH neutrino oscillations\textsuperscript{a} and the lightest $\mathcal{O}(\text{KeV})$ mass RH neutrino is a warm DM. The other two Majorana neutrinos in this scenario are much heavier but still have to be lighter than the electroweak scale as well. Previously, some other models were already proposed where the light sterile neutrino is produced by the mixing with the ordinary light ones and plays the role of Dark Matter\textsuperscript{a}. In this letter we explore the case when Majorana neutrinos mass scale is well above the electroweak scale.
where $T_{k}^{\pm}$ are destruction and production rates correspondingly (see, e.g. 12). In our case under a set of reasonable assumptions as explained in 1 this reduces to the following set of equations

$$
\dot{\rho}_{AA} = 2H_{BB}\rho_{-},
\dot{\rho}_{-} = \rho_{+}(H_{BB} - H_{AA}) - H_{AB}\rho_{BB}^{\frac{\beta}{2}},
\dot{\rho}_{+} = \rho_{-}(H_{AA} - H_{BB}),
$$

where $A, B$ are flavor indices and we have defined $\rho_{+} = (\rho_{AB} + \rho_{BA})/2, \rho_{-} = -i(\rho_{BA} - \rho_{AB})/2$. The Hamiltonian can be computed from the real part of the Majorana neutrino two-point Green function on Fig. (1). Below we will consider two cases with two different effective temperature dependent Hamiltonians result from self-energy corrections shown on Fig. (2-3).
Figure 1: On the left panel the loop induced by the interaction of the heavy CUT fermion $F$ with the Higgs and Majorana neutrinos leads to (6) with $A^{-1} \sim h^2/M_{\text{GUT}}$. Yukawa couplings are assumed to be quite small. On the right panel the tadpole diagram where $H$ is the heavy GUT Higgs, again, leads effectively to (6). Here $\mu$ is trilinear coupling between light and GUT Higgs, $y$ is the Yukawa coupling of the GUT Higgs to Majorana neutrinos. $A^{-1} \sim y\mu/M_{\text{GUT}}$.

3 The Minimal Scenario.

The Lagrangian of the model (1) is written in the Majorana neutrino mass eigenstate basis. It is more convinient to make rotation to the flavor basis where the Hamiltonian takes the form

$$H_{AB} \simeq \left( \begin{array}{cc} 0 & \frac{M^2 \Delta}{E} \\
\frac{M^2 \Delta}{E} & -\frac{M^2 \Delta}{E} + \frac{1}{2} T K_{22} \end{array} \right), \quad (3)$$

where $E = \sqrt{M^2 + k^2}$, $M = M_1 \simeq M_2$. This all is, indeed, quite similar to MSW effect with one important difference. The regime is nonadiabatic contrary to the that of the MSW effect. Assuming that $E \approx k \sim T$, $\Delta M^2_1 \equiv -2M^2 \Delta < 0$, where $\Delta \equiv (M_1 - M_2)/M_2 \ll 1$ one obtains solving the Eqs. (4)

$$\rho_{AA} \approx \frac{m_{pl}}{M} \left( \frac{K_{22}}{8\Delta} \right)^{\frac{1}{4}} K_{11} \leq \frac{m_{pl}}{2M} K_{11}, \quad (4)$$

where $m_{pl} = M_p \sqrt{90/8\pi^2 g_*} \sim 10^{18}$ GeV. The last inequality follows from the requirement that the critical temperature $T_c = M(8\Delta/K_{22})^{1/2}$ at which $H_{BB}'$ becomes equal to zero satisfies $T_c > M$. The upper bound on the DM energy density can be estimated as follows

$$\frac{\Omega_{\text{DM}} h^2}{0.1} \approx 10^{25} K_{11}, \quad (5)$$

Given that the lifetime of the flavor state $N_A$ has to exceed the age of the universe in order to be a DM candidate one can estimate that $K_{11} \leq 10^{-42}(1 \text{ TeV}/M_1)$ and the corresponding density (5) is by far too small.

4 The next-to-minimal scenario.

As it was shown in the previous section the minimal model fails to produce the correct amount of the "dark" Majorana neutrino abundance. In this section we propose an extension to the previous setup by introducing (explicitly) no new particles but allowing a dimension five effective operator

$$\mathcal{L}_{\text{eff}} = \frac{(\lambda)_{AB}}{\Lambda} |\Phi|^2 N_A^c N_B, \quad (6)$$

where $\Phi$ is usual Higgs field, $\Lambda$ is some high energy scale and the elements of the matrix $(\lambda)_{AB} \sim \mathcal{O}(1)$. The correction of the type (6) is generic if one assumes a certain additional physics at energies much higher than the electroweak scale. Some examples of the corresponding diagrams
are shown on Fig 4. Note that $\Lambda$ naturally turns out to be much larger than $M_{\text{GUT}}$ and even $M_{\text{pl}}$ scales.

The resummed Majorana propagator will contain the nondiagonal entries which are given by

$$-\Sigma_{AB} \approx \frac{1}{\Lambda} \int_0^\infty \frac{d^4q}{(2\pi)^4} \delta(q^2 - m_b^2) n_b(q),$$

where

$$n_b(q) = \frac{1}{e^{\frac{m_b}{\Lambda}} - 1},$$

with $u$ being a four velocity of the thermal bath. Neglecting the thermal mass of the Higgs one obtains the following effective Hamiltonian:

$$H \approx \begin{pmatrix} 0 & -\frac{1}{12\Lambda} & \frac{1}{12\Lambda} + \frac{1}{8} TK_{22} \\ -\frac{1}{12\Lambda} & -M_{22}^2 & TK_{22} \\ \frac{1}{12\Lambda} & TK_{22} & \frac{1}{8} \end{pmatrix}.$$ (9)

Solving Eq. (3) one obtains that

$$\rho_{AA} \sim \frac{m_{pl} M}{z_c \Lambda^2 K_{22}}.$$ (10)

where $z_c$ is given by

$$z_c = \left( \frac{K_{22}}{8\Lambda} \right)^{\frac{1}{2}}.$$ (11)

The above estimate transfers into the DM abundance which is at present

$$\Omega_{\text{DM}} h^2 = \frac{1}{0.1} \frac{\left( \frac{10^{24}\text{GeV}}{\Lambda} \right)^2 \left( \frac{M}{10^6\text{GeV}} \right)}.$$ (12)

The operator (6) after electroweak symmetry breaking leads to the mixing term $(v^2_{\text{EW}}/\Lambda) N_A N_B$ between "dark" and "bright" neutrinos. This results in the decay of the "dark" neutrino via oscillations into the "bright" one and its consequent decay. The rate of this process is

$$\Gamma(N_A \rightarrow N_B \rightarrow \ldots) = \left( \frac{v^2_{\text{EW}}}{\Lambda} \right)^2 \frac{\Gamma_B}{(\Gamma_B^f + M^2\Delta^2)}.$$ (13)

The decay rate of the "bright" neutrino is dominated by $\Gamma(N_B \rightarrow \Phi + \nu)$ and one estimate that

$$\Gamma(N_A \rightarrow N_B \rightarrow \ldots) \sim 10^{-35} z_c^4 \left( \frac{10^{24}\text{GeV}}{\Lambda} \right)^2 \left( \frac{10^6\text{GeV}}{M} \right)^2 \text{eV}.$$ (14)

There is phase space of parameters $\{M, \Lambda, z_c\}$ which leads to the correct DM abundance and the lifetime of the "dark" neutrino larger than the age of the Universe. For example, one can choose $z_c = 10^{-2}$, $M = 10^7$ GeV and $\Lambda = 10^{24}$ GeV. The "extreme" case would be to take $z_c = 10^{-3}$ (because, at smaller values it is hard to have the "bright" neutrino being thermally produced), $M = 1$ TeV and $\Lambda \sim 10^{22.5}$ GeV. In this case the "dark" neutrino has the mass near electroweak scale and a lifetime close to the present age of the Universe. Of course, a more stringent bounds may come from astrophysical observations and the lower bound on the neutrino mass scale quoted here may go up substantially.

There is another constraint which comes from the four-body decay $N_A \rightarrow 2\Phi + N_B \rightarrow 3\Phi + \nu$. The rate of such decay can be estimated as

$$\Gamma_4 \sim 10^{-7} \frac{M^3}{\Lambda^2} K_{22} \sim 10^{-24} \left( \frac{M}{10^6\text{GeV}} \right)^4 \left( \frac{10^{24}\text{GeV}}{\Lambda} \right)^2 \text{eV}.$$ (15)
Requiring that this rate is smaller then the value of the Hubble parameter at present leads to the constraint

$$\left( \frac{M}{10^3 \text{GeV}} \right)^2 \left( \frac{10^{24} \text{GeV}}{\Lambda} \right) \leq 10^{-4.5}.$$  \hspace{1cm} (16)

This implies that $M$ cannot be higher then about $10^8 - 10^9 \text{GeV}$ depending on the value of $z_c$.

5 Conclusions

In this letter we proposed the see-saw model in which one of three Majorana neutrinos ("dark" neutrino) is weakly coupled and is a DM candidate. The mechanism of the generation of a suitable abundance was suggested. At least one of the strongly coupled Majorana neutrino ("bright" neutrino) have to be quasi degenerate with the "dark" one. The allowed quasi degenerate Majoranas mass scale $M$ can, in principle, be as low as $\mathcal{O}(\text{TeV})$ and up to as high as usually suggested in GUT scenarios, i.e. $M \sim 10^8 - 10^9 \text{GeV}$. The mass of the third neutrino is adjusted correspondingly: if $M$ is below $10^8 \text{GeV}$ the third neutrino has to be also quasi degenerate with the other two in order to have lepton number produced efficiently. Otherwise it remains a relatively free parameter as in the standard hierarchical leptogenesis scenario. This letter is based on the results which will be reported in much more details in $^1$.

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LATEST RESULTS OF THE CAST AXION SEARCH


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The CAST experiment is making use of a decommissioned LHC test magnet to look for solar axions by their conversion into photons inside the magnetic field. The data taking of the first phase, with vacuum in the magnet pipes, took place in 2003 and, in improved conditions, in 2004. The final phase I result has been recently released, excluding axions down to $g_a \gamma \lesssim 8.8 \times 10^{-11}$ GeV$^{-1}$ for $m_a \lesssim 0.02$ eV. CAST is now immersed in its phase II, operating with a buffer gas inside the magnet pipes, in order to extend the sensitivity of the experiment to higher axion masses. During the 2006 data taking $^4$He was used, and now the system is being upgraded to use $^3$He. The latest status of the experiment will be presented. A brief overview of the situation of other axion experiments and in general of the field will be given.

1 Introduction

Axions are light pseudoscalar particles that arise in theories in which the Peccei-Quinn U(1) symmetry has been introduced to solve the strong CP problem\(^1\). They could have been produced in early stages of the Universe being attractive candidates to the cold Dark Matter (and in some particular scenarios to the hot Dark Matter) that could compose up to $\sim 1/3$ of the ingredients of the Universe.

Axion phenomenology\(^2\) is determined by its mass $m_a$ which in turn is fixed by the scale $f_a$ of the Peccei–Quinn symmetry breaking, $m_a \simeq 0.62$ eV ($10^7$ GeV/$f_a$). No hint is provided by theory about where the $f_a$ scale should be, so the axion mass is an unconstrained parameter on which all axion couplings depend. In addition, the particular way the axion is implemented in the Standard Model—the axion model—determines the type and magnitude of such couplings. However, only one particular process, the Primakoff effect, is present in almost every axion model and is the basis of most axion detection techniques. It makes use of the coupling between the axion field $\psi_a$ and the electromagnetic tensor:

$$\mathcal{L} = -\frac{1}{4} g_{\gamma a} \Phi^\alpha \Phi^{\bar{\alpha}} \Gamma_{\mu\nu} \Gamma^{\alpha\beta} =$$

$$= -g_{\gamma a} \psi_a \hat{B} \cdot \hat{E}$$

(1)

and allows for the conversion of the axion into a photon —and vice versa— in the presence of an electromagnetic field.

Like all other axion couplings, $g_{\gamma a}$ is proportional to $m_a^{3/4}$:

$$g_{\gamma a} \simeq 0.19 \frac{m_a}{\mathrm{eV}} \frac{E}{N} \left( \frac{1}{3} \right) \frac{2(4 + z)}{3(1 + z)} \times 10^{-9} \text{GeV}^{-1}$$

(2)

where $E/N$ is the PQ symmetry anomaly and the second term in parenthesis is the chiral symmetry breaking correction. The anomaly $E/N$ depends on the particular axion model, while the symmetry breaking correction is a function of the parameter $z = m_u/m_d \simeq 0.56$ ($m_u$ and $m_d$ being the up and down quark masses). Two popular models are the GUT–DFSZ axion\(^5\) ($E/N = 8/3$) and the KSVZ axion\(^5\) ($E/N = 0$). However, it is possible to build viable axion models with different values of $E/N$ and the determination of the parameter $z$ is subject to
some theoretical uncertainties. This implies that a very small or even vanishing $g_{\gamma\gamma}$ cannot in principle be excluded.

A combination of astrophysical and nuclear physics constraints, and the requirement that the axion relic abundance does not overclose the Universe, restricts the allowed range of viable axion masses.\cite{3,9,10} Pure cosmological arguments lead to a conservative, relative model-independent version of the allowed mass range:

$$10^{-6}\,\text{eV} \lesssim m_a \lesssim 1\,\text{eV}$$

the upper limit being recently set,\cite{11} by requiring thermal production of axions to be compatible with recent CMB data. This range and the allowed range of $g_{\gamma\gamma}$ can be further constrained by a number of theoretical arguments that depend on more or less solid astrophysical models. Let's mention the limit $g_{\gamma\gamma} \lesssim 10^{-9}\,\text{GeV}^{-1}$ based on the solar standard model and helioseismological observations,\cite{12} or the so-called globular cluster limit of $g_{\gamma\gamma} \lesssim 10^{-10}\,\text{GeV}^{-1}\,\text{eV}^{-1}\,\text{m}^{-1}\,\text{eV}^{-1}$.\cite{13,14}

Axions could be produced at early stages of the Universe by the so-called misalignment (or realignment) effect.\cite{2} Extra contributions to the relic density of non-relativistic axions might come from the decay of primordial topological defects (like axion strings or walls). There is not a consensus on how much these contributions account for, so the axion mass window may which may give the right amount of primordial axion density (to solve the dark matter problem) spans from $10^{-6}\,\text{eV}$ to $10^{-3}\,\text{eV}$. For higher masses, the axion production via these channels is normally too low to account for the missing mass, although its production via standard thermal process increases. Thermal production yields relativistic axions (hot dark matter) and is therefore less interesting from the point of view of solving the dark matter problem, but in principle axion masses up to $\sim 1\,\text{eV}$, are not in conflict with cosmological observations.\cite{11}

Under the assumption that axions are the cold dark matter, they could be detected by using microwave cavities as originally proposed in\cite{15}. In a static background magnetic field, axions will decay into single photons via the Primakoff effect. The energy of the photons is equal to the rest mass of the axion with a small contribution from its kinetic energy, hence their frequency is given by $hf = m_a c^2 (1 + O(10^{-6}))$. At the lower end of the axion mass window of interest, the frequency of the photons lies in the microwave regime. A high-Q resonant cavity, tuned to the axion mass serves as high sensitivity detector for the converted photons. Such technique is followed by experiments like the Axion Dark Matter Experiment (ADMX),\cite{16,17} which has implemented the concept using a cylindrical cavity of 50 cm in diameter and 1 m long. So far the ADMX experiment has scanned a small axion mass energy, from 1.9 to 3.3 $\mu\text{eV}$\cite{17} with a sensitivity enough to exclude a KSVZ axion, assuming that thermalized axions compose a major fraction of our galactic halo ($\rho_a = 450\,\text{MeV}/\text{c}^2$). An independent, high-resolution search channel operates in parallel to explore the possibility of fine-structure in the axion signal.\cite{18} Currently the collaboration has completed a development involving high sensitivity squids amplifiers, which will allow them to scan the full decade $10^{-6} - 10^{-5}$ eV in the near future. Plans are also ongoing both to increase the sensitivity of the experiment to lower axion-photon couplings as well as extending the reachable axion mass range.

But axions could also be copiously produced in the core of the stars by means of the Primakoff conversion of the blackbody photons in the fluctuating electric field of the plasma. In particular, a nearby and powerful source of stellar axions would be the Sun. This axion emission would open new channels of stellar energy drain. Therefore, energy loss arguments constrain considerable axion properties in order not to be in conflict with our knowledge of solar physics or stellar evolution.\cite{13}

The solar axion flux can be easily estimated\cite{19,20} within the standard solar model under the conservative assumption of an axion with no leptonic couplings (hadronic axion). The resulting

\footnote{particular scenarios with axion couplings to other particles could give rise to additional contributions to the
axion flux has an average energy of about 4 keV and can be parameterized by the following expression\textsuperscript{30}:

\[
\frac{d\Phi_a}{dE} = 6.02 \times 10^{10} \, \text{cm}^{-2} \, \text{s}^{-1} \, \text{keV}^{-1} \, g_{a\gamma}^2 \, \Theta^{1.28} \, E^{-1.20},
\]

where \( g_{a\gamma} = g_{a\gamma}/10^{-10} \, \text{GeV}^{-1} \).

By means again of the photon-axion coupling, solar axions can be converted back into photons in the presence of an electromagnetic field. The energy of the reconverted photon is equal to the incoming axion, so a flux of detectable X-rays with the same energy profile as (4) is expected after the conversion. Crystalline detectors may provide such fields\textsuperscript{21},\textsuperscript{20} giving rise to very characteristic Bragg patterns that have been looked for as byproducts of dark matter underground experiments\textsuperscript{22,23,24}. However, the prospects of this technique have been proved to be rather limited\textsuperscript{25} and do not compete with the experiments called "axion helioscopes"\textsuperscript{26,19} which use magnets to trigger the axion conversion. This technique was first experimentally applied in\textsuperscript{27} and later on by the Tokyo helioscope,\textsuperscript{28} which provided the first limit to solar axions which is "self-consistent", i.e., compatible with solar physics. Currently, the same basic concept is being used by CAST at CERN\textsuperscript{29,31,30,32,33,34} with some original additions that provide a considerable step forward in sensitivity to solar axions. In the following section we make a short description of the experiment as well as an update of its status and results.

It is worth stressing that "helioscope" experiments like CAST are not based on the assumption of axions being the dark matter. Moreover, although we focus on the axion because of its special theoretical motivations, all this scenario is also valid for a generic pseudoscalar (or scalar) particle coupled to photons\textsuperscript{35}. Needless to say that the discovery of any type of pseudoscalar or scalar fundamental particle would have profound implications in Particle Physics.

For the sake of completeness, let us mention that the existence of axions or other axion-like particles may produce measurable effects in the laboratory. A typical example is the "light through wall" experiments, in which a photon beam is converted into axions inside a magnetic field and, after crossing an optical barrier, are converted back into photons by another magnetic field. As a result, light seems to have gone through an opaque wall. This technique was used to derive some early limits on the axion properties\textsuperscript{36}.

Other subtler effects are the ones induced on the polarization of a laser beam traversing a magnetic field in vacuum. The presence of axion-photon oscillations will produce both a rotation (dichroism) and an ellipticity of the beam polarization. Although the ellipticity effect has a Standard Model contribution, by virtue of four-legged fermion loops, the dichroism one does not. Experiments with ultra-precise optical equipment may look for such an effect. The PVLAS experiment\textsuperscript{39}, designed to measure the QED-predicted magnetic-induced birefringence\textsuperscript{41,40} has recently reported on a positive detection\textsuperscript{42} compatible in principle with the presence of a photonic oscillation\textsuperscript{3}. However, the interpretation of PVLAS observation in terms of axions needs an axion mass of \( \sim 1 \, \text{meV} \) and an axion-photon coupling of \( \sim 10^{-6} \, \text{GeV}^{-1} \), far larger than many astrophysical limits and experimental results, in particular that of CAST (although exotic extensions of the standard axion scenario may allow to reconcile all experimental results\textsuperscript{43}). The current situation of these type of experiments and theoretical efforts was reviewed in this same conference by J. Jaeckel, and we refer to his contribution for further information and references.
2 The CAST experiment

The CAST experiment is making use of a decommissioned LHC test magnet that provides a magnetic field of 9 Tesla along its two parallel pipes of $2 \times 14.5$ cm$^2$ area and 10 m length. The aperture of each of the bores fully covers the potentially axion-emitting solar core ($\sim 1/10$th of the solar radius). The magnet is mounted on a platform with $\pm 8^\circ$ vertical movement, allowing for observation of the Sun for 1.5 h at both sunrise and sunset. The rest of the day is devoted to background measurements. The horizontal range of $\pm 40^\circ$ encompasses nearly the full azimuthal movement of the Sun throughout the year. At both ends of the magnet, several detectors look for the X-rays originated by the conversion of the axions inside the magnet when it is pointing to the Sun.

These features makes the axion-photon conversion probability in the CAST magnet be a factor 100 higher than in the previous best helioscope at Tokyo. More specifically, the probability that an axion going through the transverse magnetic field $B$ over a length $L$ will convert to a photon is given by:

$$P_{\gamma} = 2.4 \times 10^{-17} \left( \frac{B}{9.6 \text{ T}} \right)^2 \left( \frac{L}{10 \text{ m}} \right)^2 \left( g_{\gamma} \times 10^{10} \text{ GeV}^{-1} \right) |M|^2$$

where the matrix element $|M|^2$ accounts for the coherence of the process:

$$|M|^2 = 2(1 - \cos qL)/(qL)^2$$

being $q$ the momentum exchange. The fact that the axion is not massless, makes the axion and photon waves out of phase after a certain length. For axion energies relevant to us and the length of the magnet, the coherence is preserved ($|M|^2 \simeq 1$) for axion masses up to $\sim 10^{-2}$ eV, while for higher masses $|M|$ begins to decrease, and so does the sensitivity of the experiment. To cope with this, a second phase of CAST was planned with the magnet beam pipes filled with a buffer gas to give a mass to the photons $m_\gamma = \omega_p$ (where $\omega_p$ is the plasma frequency of the gas, $\omega_p^2 = 4\pi\epsilon_0 n_e \rho_0$), being $n_e$ the spatial density of electrons and $\rho_0$ the classical electron radius). For axion masses that match the photon mass, the coherence is restored. Changing the pressure of the gas inside the pipe, the photon mass can be changed accordingly, and so the sensitivity of the experiment can be extended to higher axion masses.

A full cryogenic station is used to cool the superconducting magnet down to 1.8 K$^{44}$ The hardware and software of the tracking system have been precisely calibrated, by means of geometric survey measurements, in order to orient the magnet to any given celestial coordinates. The overall CAST pointing precision is better than 0.01$^\circ$.45

At both ends of the magnet, three different detectors search for excess X-rays from axion conversion in the magnet when it is pointing to the Sun. Covering both bores of one of the magnet’s ends, a conventional Time Projection Chamber (TPC) is looking for X-rays from “sunrise” axions. At the other end, facing “sunset” axions, a second smaller gaseous chamber with novel MICROMEGAS (micromesh gaseous structure – MM) $^{46}$ readout is placed behind one of the magnet bores, while in the other one a focusing X-ray mirror telescope is working with a Charge Coupled Device (CCD) as the focal plane detector. Both the CCD and the X-ray telescope are prototypes developed for X-ray astronomy. $^{47}$ The X-ray mirror telescope can produce an “axion image” of the Sun by focusing the photons from axion conversion to a $\sim 6 \text{ mm}^2$ spot on the CCD. The enhanced signal-to-background ratio substantially improves the sensitivity of the experiment.
CAST phase II: status

CAST has been running in phase I (vacuum) configuration both in 2003 and 2004. The combined results of these data for all three detectors have been recently released.\textsuperscript{30} From the absence of signal in the data a 95% CL upper limit to the axion-photon coupling was derived:

\[ g_{a\gamma} < 8.8 \times 10^{-11}\text{GeV}^{-1}(95\%\text{C.L.}) \]  

(7)

This limit is valid for the mass range \( m_a \lesssim 0.02 \text{ eV} \) where the expected signal is mass-independent because, as was explained in the introduction, the axion-photon oscillation length far exceeds the length of the magnet and hence coherence is preserved in the conversion. For higher \( m_a \) the overall signal strength diminishes rapidly and the spectral shape differs. The limit to the axion-photon coupling was derived also for axion masses above this range so that the entire 95% CL exclusion line shown in Fig. 1 was obtained. This result improves the previous CAST published result\textsuperscript{31} and constitutes the final official CAST phase I limit. As can be seen, it is a factor \( \sim 7 \) more restrictive than the limit from the Tokyo axion helioscope and goes for the first time beyond the limit derived from stellar energy-loss arguments.

During 2005 the experiment was upgraded to face the needs of phase II operation, which require the injection of a buffer gas in the magnet pipes and the precise control of its pressure. As a result, a system dealing with \(^4\text{He}\) gas for that purpose was built and has been operational since end of 2005. Data taking with \(^4\text{He}\) gas in the magnet pipes took place all throughout 2006,
setting a different pressure every day, so the axion mass range was scanned continuously up to the condensation limit of this gas at the operation temperature of 1.8 K, which is \(~13\) mbar. This has allowed CAST to scan a new axion mass range up to \(~0.39\) eV, exploring the region indicated in fig. 2. The \(^4\)He data is currently under analysis, the region shown in figure 2 being the expected exclusion plot.

Currently the experiment is being further upgraded to be able to deal with \(^3\)He as the buffer gas. This gas can reach pressures up to \(~135\) mbar at 1.8 K, corresponding to axion masses up to \(~1.2\) eV. The experimental setup to use \(^3\)He, supposes a considerable effort due to the high safety and reliability level required by the fact of not tolerating leaks of this gas, considerable expensive. The system is expected to be ready for data taking by mid 2007.

For the realization of the \(^3\)He experimental program and to exploit fully the potential of the new upgraded system, CAST plans to run during the next 3 years, scanning axion masses up to \(~1.2\) eV, and closing the window presently allowed (axions with masses above \(1.1\) eV are excluded by the amount of Hot Dark Matter induced by the cosmic microwave background data). The expected reachable region is depicted in figure 2. As can be seen, the second phase of CAST is allowing sensitivity to QCD axion models at the \(0.1 - 1\) eV mass scale, which was out of reach for previous experiments (including CAST phase I).

In parallel with the above mentioned upgrades, CAST is exploring possible improvements in the detector systems that could lead to an increased overall sensitivity. A very appealing possibility is to add a second X-ray focusing optics to the experimental setup. A design has been done specifically for CAST, and tests are ongoing to assess whether construction under specification is possible. It consists of a concentrator with a 1.3 m focal length and 47 mm diameter built using new substrate techniques developed at LLNL of Livermore. The concentrator will have 14 nested polycarbonate conic shells, each 125 mm long and coated with iridium. It will be coupled to a new smaller Micromegas detector with enhanced features with respect to
the present version, in particular it will enjoy a shielding composed of copper, lead, cadmium nitrogen and polyethylene following the experience of the TPC detector described above, and is expected to reduce the background similarly.\textsuperscript{25}

4 Conclusions

The CAST experiment is looking for solar axions following the "axion helioscope" concept with a 9.6 Tesla and 10 m long LHC test magnet. A final result from the phase I data taken in 2003 and 2004 data has been presented: $g_{ax} < 0.88 \times 10^{-11}$ GeV\(^{-1}\) for $m_a \lesssim 0.02$ eV. The phase II of the experiment has already started using \(^4\)He as buffer gas to trigger axion-photon conversion for higher axion masses. With this upgrade CAST has entered the theory motivated axion parameter space. Currently the collaboration prepares the transition to \(^3\)He which will allow to extend the sensitivity of the experiment to axion masses up to $\sim 1.2$ eV, closing the window allowed by cosmological limits. The \(^3\)He phase should start in the coming months and last for about three years.

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Neutralino dark matter in the NMSSM: phenomenological viability

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We study the viability of the lightest neutralino as a dark matter candidate in the Next-to-Minimal Supersymmetric Standard Model (NMSSM). In our analysis we take into account accelerator constraints as well as bounds on low-energy observables (muon anomalous magnetic moment, rare K and B meson decays). We further impose consistency with present bounds on the neutralino relic density. We also address the prospects for the direct detection of neutralino dark matter in the allowed regions of the parameter space, comparing the results with the sensitivities of present and projected dark matter experiments. We find regions of the NMSSM parameter space where the neutralino has the correct relic abundance and its detection cross section is within the reach of dark matter detectors, essentially owing to the presence of very light singlet-like Higgses, and either singlino dominated or very light neutralinos.

1 Introduction

Supersymmetric (SUSY) models with R-parity conservation offer excellent candidates for dark matter. In particular, the lightest neutralino ($\chi_1^0$) is one of the most interesting within the class of Weakly Interactive Massive Particles (WIMPs). WIMPs can in principle be directly detected via elastic scattering on target nuclei, and there are currently a large number of experiments devoted to the direct detection of WIMP dark matter $^1$.

The Next-to-Minimal Supersymmetric Standard Model (NMSSM) is an extension of the Minimal Supersymmetric Standard Model (MSSM) by a singlet superfield $\hat{S}$. The NMSSM provides an elegant solution to the so-called $\mu$ problem of the MSSM, while at the same time rendering the Higgs “little fine tuning problem” of the MSSM less severe. The presence of additional fields, namely an extra CP-even and CP-odd neutral Higgs bosons, as well as a fifth neutralino, leads to a richer and more complex phenomenology. This also translates into the possibility of dark matter scenarios that can be very different from those encountered in the MSSM, both regarding the relic density and the prospects for direct detection. In particular, the
exchange of very light Higgses can lead to large direct detection cross sections, within the reach of the present generation of dark matter detectors\(^2\). A systematic analysis of the low-energy NMSSM parameter space has recently been conducted\(^3\), including the constraints from LEPII and Tevatron as well as those from the SUSY contribution to the muon anomalous magnetic moment, and bounds from \(K\)- and \(B\)-meson decays. By further including the constraints on the neutralino relic density, we evaluate the prospects for the neutralino detection cross section on the allowed regions of the parameter space, comparing the results with the sensitivity of dark matter detectors.

2 Constraints on the NMSSM low-energy parameter space

The addition of a gauge singlet superfield \(\tilde{S}\) modifies the MSSM superpotential as follows:

\[
W_{\text{NMSSM}} = \epsilon_{ij} \left( Y_u \hat{H}^c_i \hat{Q}^j \hat{u} + Y_d \hat{H}^c_i \hat{Q}^j \hat{d} + Y_e \hat{H}^c_i \hat{L}^j \hat{e} \right) - \epsilon_{ij} \lambda \hat{S} \hat{H}^c_i \hat{H}^c_j + \frac{1}{3} \kappa \hat{S}^3. \tag{1}
\]

After the spontaneous breaking of electroweak (EW) symmetry, the neutral Higgs scalars develop vacuum expectation values (VEVs), \(\langle H_1^0 \rangle = v_1\), \(\langle H_2^0 \rangle = v_2\) and \(\langle S \rangle = s\). This leads to the dynamical generation of an effective interaction \(\mu \hat{H}_1 \hat{H}_2\), with \(\mu \equiv \lambda s\).

In the NMSSM spectrum, we now have three CP-even and two CP-odd Higgs states. In particular, the lightest Higgs scalar can be written as \(h_1^0 = S_1 \hat{H}_1^0 + S_2 \hat{H}_2^0 + S_3 \hat{S}\), where \(S\) is the unitary matrix that diagonalises the \(3 \times 3\) scalar Higgs mass matrix. In the neutralino sector, the singlinos mix with the bino, wino and Higgsinos. The lightest state can be now expressed as \(\tilde{\chi}_1^0 = N_{11} \tilde{B}^0 + N_{12} \tilde{W}_3^0 + N_{13} \tilde{H}_1^0 + N_{14} \tilde{H}_2^0 + N_{15} \tilde{S}\), where \(N\) diagonalises the \(5 \times 5\) neutralino mass matrix.

The low-energy NMSSM parameter space can be described in terms of \(\lambda, \kappa, \tan \beta, \mu, A_\lambda, \)\(\) \(A_\kappa\) degrees of freedom, as well as the soft SUSY-breaking terms, namely gaugino masses, \(m_{Q, L, U, D, E}\), scalar masses, \(m_{\phi, L, U, D, E}\), and trilinear parameters, \(A_{Q, L, U, D, E}\). A thorough analysis of the low-energy NMSSM phenomenology (minimisation of the potential, computation of spectrum and compatibility with LEP/Tevatron bounds) can be obtained using the NMHDECAY 2.0 code\(^4\). Additionally, we have also included in our analysis\(^3\) a more precise computation of the \(b \rightarrow s\gamma\) decay in the NMSSM\(^5\), taking into account next-to-leading order contributions, and imposing consistency at the \(2\sigma\) level with the experimental central value\(^6\),

\[
\text{BR}^{\text{exp}}(b \rightarrow s\gamma) = (3.55 \pm 0.27 \times 10^{-4}). \tag{2}
\]

Likewise, we have also incorporated the constraints coming from the contribution of a light pseudoscalar \(a^0\) in the NMSSM to the rare \(B\)- and \(K\)-meson decays\(^7\). Finally, in our analysis we have also included the constraints coming from the SUSY contributions to the muon anomalous magnetic moment, \(a_\mu = (g_\mu - 2)\). At present, the observed excess in \(a_\mu^{\text{exp}}\) constrains a possible SUSY contribution to be \(8.6^{\text{SUSY}} = (27.6 \pm 8) \times 10^{-10}\). Concerning the evaluation of the SUSY contributions to \(a_\mu\), the only change with respect to the MSSM is due to the fifth neutralino state and the corresponding modified neutralino-lepton-slepton coupling. For the regions of the parameter space exhibiting good prospects regarding the direct detection of dark matter\(^2\), the SUSY contributions are in general quite small. A sufficiently large \(a_\mu^{\text{SUSY}}\) can nevertheless be obtained when slepton (and gaugino) masses are decreased, in association with large values of the slepton trilinear coupling\(^5\).

Regarding the bounds arising from \(K\)- and \(B\)-meson physics, the most important role is played by the \(b \rightarrow s\gamma\) decay, which can in principle exclude important regions of the parameter

\(^{a}\)We impose at low energies a relation for \(M_1\), that mimics a hypothetical GUT unification, \(M_3 = 2M_2 = 6M_1\).

\(^{b}\)For example, assuming \(m_E = 150\) GeV and \(A_E = -2500\) GeV, and setting bino mass to \(M_1 = 100\) GeV, leads to a sufficiently large \(a_\mu^{\text{SUSY}}(\mathcal{O}(10^{-9}))\).
space. Under our assumptions $^c$, the most important contributions to \(\text{BR}(b \to s\gamma)\) arise in general from charged Higgs diagrams$^3$. In the NMSSM, the charged Higgs masses are given by

$$m_{H^\pm}^2 = \frac{2\mu^2}{\sin(2\beta)} \frac{\kappa}{\lambda} - v^2 \lambda^2 + \frac{2\mu A_\lambda}{\sin(2\beta)} + M_1^2,$$  \hspace{1cm} (3)

leading to the conclusion that larger values of \(m_{H^\pm}^2\) and thus smaller \(\text{BR}(b \to s\gamma)\), should be obtained when \(\kappa/\lambda\) is sizable (for positive values of \(\kappa\)) or for small \(\kappa/\lambda\) (if \(\kappa < 0\)). In general, smaller values of the \(\text{BR}(b \to s\gamma)\) will be also associated to larger values of the product \(\mu A_\lambda\) and to larger values of tan\(\beta\).

As an example $^d$, we represent in Fig. 1 the \((\lambda, \kappa)\) parameter space for tan\(\beta = 3\), \(A_\lambda = -200\text{ GeV}\), \(A_\kappa = -200\text{ GeV}\) and \(\mu = 130\text{ GeV}\). Exclusion areas due to the violation of theoretical (Landau poles, false minima, tachyons) and/or experimental constraints (in this case due to conflict with LEP/Tevatron data) are depicted. The isosurfaces for \(\text{BR}(b \to s\gamma)\) on the \((\lambda, \kappa)\) plane are also displayed. The resulting branching ratio is typically large, especially in regions with small \(\kappa/\lambda\), where the charged Higgs mass is smaller. In this example, only a small triangular region with \(\lambda \lesssim 0.05\), for \(\kappa < 0.7\), is within a 1\sigma deviation from the experimental bound of Eq. (2) and \(\lambda \lesssim 0.35\) is needed in order to be within 2\sigma of that result. In the plot we also indicate with dot-dashed lines the different values of the charged Higgs mass, thus illustrating the correlation between its decrease and the increase in \(\text{BR}(b \to s\gamma)\).

In order to be a good dark matter candidate, the lightest NMSSM neutralino must also comply with the increasingly stringent bounds on its relic density. Astrophysical constraints$^1$ suggest the following range for the WIMP relic abundance

$$0.1 \lesssim \Omega h^2 \lesssim 0.3,$$ \hspace{1cm} (4)

which can be further reduced to

$$0.095 \lesssim \Omega h^2 \lesssim 0.112,$$ \hspace{1cm} (5)

$^c$We do not take into account any source of flavour violation other than the Cabibbo-Kobayashi-Maskawa matrix. We will also be systematically considering large values for the squark and gluino masses (above 1 TeV).

$^d$For a comprehensive study of the parameter space see$^3$. 
taking into account the recent three years data from the WMAP satellite. Compared to what occurs in the MSSM, one would expect several alterations regarding the dominant processes that lead to $\Omega_{\chi^0_1} h^2$: first, and given the presence of a fifth neutralino (singlino), the composition of the annihilating WIMPs can be significantly different. The possibility of a singlino-like lightest supersymmetric particle (LSP), associated with new couplings in the interaction Lagrangian, may favour the coupling of WIMPs to a singlet-like Higgs, whose mass can be substantially lighter than in the MSSM, given the more relaxed experimental constraints. Secondly, in the NMSSM we have new open channels for neutralino annihilation. For instance, the presence of additional Higgs states may favour annihilation via $s$-channel resonances. On the other hand, light h_1^0 and a_1^0 states, that are experimentally viable, suggest that new channels with annihilation into $Z h_1^0$, $h_1^0 h_1^0$, $h_1^0 a_1^0$ and $a_1^0 a_1^0$ (either via $s$-channel $Z$, $h_1^0$, $a_1^0$ exchange or $t$-channel neutralino exchange) can provide important contributions to the annihilation and co-annihilation cross-sections.

Since the goal of our work was to discuss the potential of NMSSM-like scenarios regarding the theoretical predictions for $\sigma_{\chi^0_1-p}$, we focus on the regions of the parameter space likely to have large neutralino detection cross sections. As an example, let us take $M_1 = 160$ GeV, $A_\lambda = 400$ GeV, $A_\kappa = -200$ GeV, and $\mu = 130$ GeV, with $\tan \beta = 5$, which is consistent with bounds on $\sigma_{\mu}^{\text{SUSY}}$ and $\text{BR}(b \to s \gamma)$. The results for the neutralino relic density, obtained from an NMHDECAY link to MicrOMEGAS, are depicted in the $(\lambda, \kappa)$ plane on Fig. 2.

For large values of $\kappa$ and small $\lambda$, the lightest neutralino is relatively heavy and has a mixed bino-Higgsino composition. Due to its important Higgsino component, the relic density is too small to account for $\Omega_{\chi^0_1} h^2$. Moving towards smaller values of $\kappa$ and larger values of $\lambda$, the neutralino becomes lighter and has a larger singlino component, thus leading to an increase in $\Omega_{\chi^0_1} h^2$. As the neutralino mass decreases, the annihilation channels become kinematically forbidden, such as annihilation into a pair of $Z$ or $W$ bosons when $m_{\chi^0_1} < M_Z$ or $m_{\chi^0_1} < M_W$, respectively. Below these, the resulting relic density can be large enough to fulfil the WMAP constraint. Notice that the mass and composition of the lightest Higgs can also play a key role, given that when the Higgs is sufficiently light new annihilation channels are available for the neutralino, thus decreasing its relic density.
As we can see, in the present example the correct relic density is only obtained when either the singlino composition of the neutralino is large enough or when the annihilation channels into $Z$, $W$, or $h^0$ are kinematically forbidden. Interestingly, some allowed areas are very close to the tachyonic border, which as we will verify, can give rise to very large direct detection cross sections.

3 Prospects for NMSSM direct dark matter detection

As pointed out in the literature, the existence of a fifth neutralino state, together with the presence of new terms in the Higgs-neutralino-neutralino interaction (which are proportional to $\lambda$ and $\kappa$), trigger new contributions to the spin-independent part of the neutralino-nucleon cross section, $\sigma^{\text{SI}}_{\chi_1^0-p}$. On the one hand, although the term associated with the $s$-channel squark exchange is formally identical to the MSSM case, it can be significantly reduced if the lightest neutralino has a major singlino composition. On the other hand, and more importantly, the dominant contribution to $\sigma^{\text{SI}}_{\chi_1^0-p}$ associated to the exchange of CP-even Higgs bosons on the $t$-channel, can be largely enhanced when these are very light. Consequently, large detection cross sections can be obtained, even within the reach of the present generation of dark matter detectors. Let us begin by revisiting the same example as in Fig. 2, displayed on the left-hand side of Fig. 3.

Regions of the parameter space where the neutralino fulfills all experimental constraints and has the correct relic density can be found. The latter are characterized by neutralinos with a significant singlino fraction and/or a small mass. In this case, one of the allowed regions is close to the tachyonic area and exhibits very light singlet-like Higgses, potentially leading to large detection cross sections. This is indeed the case, as evidenced on the left-hand side of Fig. 3, where the theoretical predictions for $\sigma^{\text{SI}}_{\chi_1^0-p}$ are plotted versus the lightest neutralino mass. The resulting $\sigma^{\text{SI}}_{\chi_1^0-p}$ spans several orders of magnitude, but, remarkably, areas with $\sigma^{\text{SI}}_{\chi_1^0-p} \gtrsim 10^{-7}$ pb are found. These correspond to the above mentioned regions of the parameter space with very light singlet-like Higgses ($25 \text{ GeV} \lesssim m_{h^0} \lesssim 50 \text{ GeV}$ with $S_{13}^2 \gtrsim 0.99$). The neutralino is a mixed singlino-Higgsino state ($N_{15}^2 \approx 0.35$) with mass around 75 GeV. The sensitivities of present and projected dark matter experiments are also depicted for comparison.

On the right-hand side of Fig. 3 we show the resulting $\sigma^{\text{SI}}_{\chi_1^0-p}$ when the neutralino composition is changed, namely when the Higgsino component is enhanced. Such neutralinos annihilate more efficiently, thus leading to a reduced $\Omega_{\chi_1^0} h^2$, so that the astrophysical constraint becomes more stringent. On the right-hand side of Fig. 3, the various resonances appear as funnels in the predicted $\sigma^{\text{SI}}_{\chi_1^0-p}$ for the regions with the correct $\Omega_{\chi_1^0} h^2$ at the corresponding values of the neutralino mass ($m_{\chi_1^0} \approx M_Z/2$ and $m_{\chi_1^0} \approx m_{\mu/2}$). Below the resonance with the $Z$ boson, light neutralinos are obtained $m_{\chi_1^0} \lesssim M_Z/2$ with a large singlino composition which have the correct relic abundance. The lightest Higgs is also singlet-like and very light, leading to a very large detection cross section, $\sigma^{\text{SI}}_{\chi_1^0-p} \gtrsim 10^{-6}$ pb.

4 Conclusions

We have carried out a systematic analysis of the low-energy parameter space of the Next-to-Minimal Supersymmetric Standard Model (NMSSM), addressing the implications of experimental and astrophysical constraints on the direct detection of neutralino dark matter. We have found very stringent constraints on the parameter space coming from low-energy observables, especially $\alpha_{\mu}^\text{SUSY}$ and $b \rightarrow s \gamma$. Compatibility with the neutralino relic density leads us to regions of the parameter space where either the neutralino mass is small enough for some annihilation channels to be kinematically forbidden or when the singlino component of the lightest neutralino is large enough to suppress neutralino annihilation. Some of the regions fulfilling all the experimental
and astrophysical constraints display very light, singlet-like Higgses, and are associated with very large values of $\sigma_{\tilde{\chi}}^0 - p_T$, even within the reach of dark matter experiments. In addition, the presence of singlino-Higgsino-like neutralinos is also representative of the NMSSM.

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SUPERSYMMETRIC LARGE EXTRA DIMENSIONS

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This article reviews the arguments why extra dimensions provide a unique opportunity for progress on the cosmological constant problem, and updates the status of—and the objections to (with replies)—the specific proposal using supersymmetric large extra dimensions (SLED).

1 Extra Dimensions and the Cosmological Constant Problem

For thirty years technical naturalness—the requirement that small parameters be stable against renormalization—has been a major guideline for searches to replace the Standard Model. For instance, the observation that particles with mass $M$ contribute to the Higgs potential an amount $\delta V_H \propto M^2 H^* H$ leads to the Hierarchy Problem: how can $M_W/M_p \sim 10^{-15}$ be technically natural if any particles at all have masses between $M_W$ and $M_p$? Naturalness would be assured if the Higgs were composite at a scale $\Lambda_c \gtrsim M_W$ since then there is no potential for heavy particles to correct above the scale $\Lambda_c$. It would also be assured if the Higgs were elementary but supersymmetry, broken at scales $\Lambda_s \gtrsim M_s$, enforced cancellation of bosons and fermions in their contribution to $\delta V_H$. Such considerations significantly shaped the design of the LHC, in order to test both of these proposals.

Naturalness in crisis

Yet the discovery that the universe is now entering an epoch of accelerated expansion has provoked an unprecedented angst about the use of technical naturalness as a fundamental theoretical criterion. It does so because the acceleration is well described by adding a cosmological constant, $\lambda$, to Einstein's equations,

$$R_{\mu\nu} - \frac{1}{2} R g_{\mu\nu} + \lambda g_{\mu\nu} = \frac{T_{\mu\nu}}{M_p^2},$$

(1)
with the required constant much smaller than most other fundamental scales we know. Regarded as an energy density, \( \lambda \equiv \rho/M_p^4 \), observations require \( \rho = \mu^4 \), with \( \mu \lesssim 10^{-2} \text{ eV} \).

Since particles of mass \( M \) contribute \( \delta \rho \propto M^4 \), the contribution of almost all known particles to \( \rho \) are already too much too large: for electrons \( m_e^4/\mu^4 \sim 10^{36} \), for the QCD phase transition \( A_{QCD}^4/\mu^4 \sim 10^{44} \), while for electroweak bosons \( M_W^4/\mu^4 \sim 10^{66} \). The contributions of particles with \( M \gg M_w \) are generically larger still, but can be suppressed (such as by supersymmetry) to contribute only \( \delta \rho \sim M^4_\lambda \).

This makes the cosmological constant (CC) problem the mother of all naturalness problems, since its roots lie with particles we already know rather than hypothetical particles having \( M \gg M_w \). With naturalness as their guide, theorists unaware of accelerators operating above \( 10^{-2} \text{ eV} \) could confidently predict the discovery of supersymmetric partners split in mass from the electron by this scale, in order to solve the CC problem. How can we trust naturalness as a guide at the electroweak scale if it lets us down so badly on scales we thought we understood?

**How extra dimensions can help**

The essence of the problem is that the Lorentz invariance of the vacuum implies that the vacuum stress energy satisfies \( T_{\mu \nu}^{\text{vac}} = -\rho g_{\mu \nu} \), making \( \rho/M_p^2 \) indistinguishable from \( \lambda \) in eq. (1). The puzzle is how the curvature of space (and so also \( \lambda \)) can be as small as is measured when quantum corrections to \( \rho \) should be large.

Extra dimensions help by breaking the link between 4D Lorentz invariant energies (\( \rho \)) and 4D curvature (\( \lambda \)). They can do so because if the vacuum energy is associated with the tension of a surface (or brane), then it is localized (and not Lorentz invariant) in the extra dimensions. Although it necessarily curves spacetime, it sometimes does so by curving the extra dimensions and not the four dimensions we see.\(^5,6,7\)

This is all very well, but any extra dimensional model becomes effectively four dimensional at energies above its Kaluza-Klein scale, \( \Lambda_K \sim 1/r \), where \( r \) denotes a generic linear size (radius) for the largest extra dimensions. Consequently an intrinsically extra-dimensional explanation of the size of \( \rho \) can only be useful if the extra dimensions are large: \( \Lambda_K \) cannot be too much larger than \( \mu \sim 10^{-2} \text{ eV} \), so \( r \) can’t be much smaller than \( \sim 10 \text{ m} \). Remarkably, extra dimensions can actually be this large,\(^8\) but only within a ‘brane-world’ scenario for which all observed particles are trapped on a 4D surface (or 3-brane). In this case only gravitational measurements probe the extra dimensions, and constraints on deviations from Newton’s Law presently allow dimensions slightly smaller than \( 50 \text{ m} \).\(^9\) Most encouragingly, large extra dimensions potentially do just what one wants: because observed particles are trapped on a brane, their non-gravitational properties are unchanged (as they must be) at the energies to which we have access. All that is modified is how their vacuum energy gravitates.

The extra-dimensional approach to the CC problem starts with this observation, and asks whether the theoretical elbow room thus opened is sufficient to allow a small enough 4D curvature in a technically natural way. This involves re-asking the cosmological constant problem in higher dimensions: What choices are required to make our observed 4 dimensions very flat? And can these choices be stable against renormalization? So far these issues are most thoroughly explored in 6 dimensions, to which we now turn.

2 Supersymmetric Large Extra Dimensions (SLED)

The best-developed proposal along these lines is the SLED proposal,\(^7,10,11,12\) according to which all known particles are localized on one of possibly many (usually two) parallel 3-branes that are situated at points within a 6D spacetime whose two extra dimensions are at present imagined to be \( r \sim 10 \text{ m} \) in size, so that \( 1/r \sim 10^{-2} \text{ eV} \) is not so different from \( \mu \). It is further
assumed that the ‘bulk’ physics – not trapped on the branes – is supersymmetric, and so is described by any one of the many known 6D supergravities. If the extra dimensions are not too strongly warped (as is true for the majority of explicit solutions known\textsuperscript{13,10}) the 4D Planck mass is of order\textsuperscript{8} \( M_p \sim M_g^2 r \), so the scale of the 6D Newton constant must be \( M_g \sim 10 \) TeV. The bulk supersymmetry is imagined to be badly broken by the branes, whose tension is imagined to be of the order of (but somewhat smaller than) \( M_g \).

This proposal is the best developed in several senses. First, it is the one for which the naturalness issues have been the most thoroughly explored.\textsuperscript{14,15,16,17,18} Second, it is (so far) the only extra-dimensional framework that does not argue for a vanishing 4D curvature, but instead provides an explicit mechanism for a nonzero 4D curvature of size \( \mu \). Finally, it leads to a known low-energy 4D field theory within which gravity is described by a scalar-tensor system, with the scalar labelling the classical flat direction corresponding to overall re-scalings of the extra dimensions, within which a realistic quintessence-type accelerated expansion can plausibly take place.\textsuperscript{19,20,21}

Best of all, because the extra dimensions themselves must be very large, \( 1/r \sim \mu \), and the scale of gravity in the extra dimensions must be low, \( M_g \sim 10 \) TeV. As a result the proposal is unusually predictive — with many testable predictions for tests of gravity and for particle colliders,\textsuperscript{22} in addition to its implications for cosmology.

2.1 Where we stand

The SLED proposal involves re-asking the CC problem in higher dimensions. This comes in two steps: (i) enumerate the choices which are required to obtain a small 4D curvature within a particular extra-dimensional context; (ii) identify whether or not these choices are stable against renormalization (and so are technically natural).

What is required for 4D flatness?

Most of the progress so far has been in identifying what choices are required for matter on the branes in order to obtain compactifications whose 4D geometry is approximately flat. Although it is not crucial for the naturalness arguments, these choices are best explored within chiral gauged 6D supergravity.\textsuperscript{23} Although not the simplest, this supergravity receives special attention because it allows spherical compactifications, and so can involve only positive-tension branes.\textsuperscript{24,7,10,13} Spherical extra dimensions are related to positive-tension branes by a topological argument, which is easiest to see for branes whose tension, \( T_b \), back-reacts on the geometry to give a conical defect, with defect angle \( \delta_b = T_b/M_g^4 \) (so \( \delta_b/2\pi = 4GT_b^2 \)). The Euler number, \( \chi \), for the extra dimensions then is

\[
\chi = 4G \sum_b T_b + \frac{1}{4\pi} \int d^2 x \sqrt{g} R_2 ,
\]

where \( R_2 \) is the 2D geometry's Ricci scalar. Notice that for toroidal compactifications \( R_2 = \chi = 0 \), so the brane tensions must all sum to zero (and some in particular must be negative). On the other hand, for spherical compactifications \( \chi = 2 \), so all of the tensions can be positive.

A broad class of exact solutions to 6D gauged, chiral supergravity are now known, including those which are 4D flat\textsuperscript{7,10,13,18} those having curved 4D maximal symmetry\textsuperscript{15} and those which are time-dependent.\textsuperscript{16} These show that solutions appropriate to two source branes are generically time-dependent, describing geometries wherein the extra dimensions implode or run away to flat 6D space. It turns out that for codimension-two branes a sufficient condition that ensures that all static solutions are 4D flat is to have the branes not couple to the 6D dilaton (a scalar which is partnered to the graviton by 6D supersymmetry).\textsuperscript{10,16}
How stable are these choices?

The next question asks how stable are the choices required to make the observed 4 dimensions flat. This question comes in two separate parts: (i) are the choices required for couplings in the action stable against renormalization; and (ii) given specific choices for the action, are acceptable solutions stable against changes to the initial conditions.

**Stability to initial conditions:** Given the number of solutions now known it is clear that, even given appropriate brane properties, the generic solutions to 6D supergravity describe time-dependent runaways.\textsuperscript{16} This shows that extra dimensional approaches to the CC problem generically have an initial condition problem: they describe the universe around us only if the universe starts out in a particular way. This makes them like the Hot Big Bang model itself, whose similar initial-condition problems inspired the invention of inflationary scenarios. Since the plausibility of initial conditions for the later universe can potentially be addressed by changing the dynamics of the earlier universe (such as through inflation), this kind of initial-condition problem is a price worth paying if it allows progress to be made on the more difficult issue of technical naturalness.

**Stability to renormalization:** The key question is whether the choices which allow 4D flat solutions are natural, in the sense of being stable against renormalization. Once arranged as desired, do these choices stay made as heavy particles are integrated out? Although work along these lines is still in progress, some partial results are known.\textsuperscript{14}

It is known that the Casimir energy produced by integrating out bulk fields for a toroidal bulk have the desired size. More generally, integrating out heavy bulk particles at one loop tends not to cause problems because these loops know about the full 6D supersymmetry of the action. It is the 6D supersymmetry which is relevant to integrating out the dangerous frequencies, $\omega \gtrsim M_{6}$, even though 6D SUSY is broken by the background geometry. This is because these dangerous modes probe very short distances in the bulk, and so are largely insensitive to the geometry over scales $\sim r$ (as they would have to be to ‘know’ that supersymmetry breaks).

In some circumstances integrating out massive brane fields also need not be dangerous, despite supersymmetry being badly broken on the branes. This is because a sufficient condition for 4D flatness is the absence of a brane coupling to the bulk dilaton, and arbitrary numbers of brane loops cannot generate a coupling to the dilaton if it is not already present at the classical level.

The potentially most dangerous contributions are those which mix brane and bulk loops, since these can introduce couplings between the brane and the dilaton and can know about supersymmetry breaking. The good news here is that it is sufficient to establish that these contributions are small to a small number of loops in the bulk, because the very large size of the extra dimension implies the bulk loops cost a factor of order $1/r^{4}$. (Recall that the observations require a 4D energy density of order $\mu^{4} \sim 1/r^{4}$.) It is these calculations of naturalness on which the success or failure of the SLED proposal must ultimately be judged.

3 Some Objections

It is useful to close by listing some of the best objections which have been raised against the SLED proposal over the years, together with a cartoon of the arguments as to why they do not (yet) appear to be show-stoppers.

**Why isn’t SLED killed by the arguments against 5D self-tuning?**

The observation that higher dimensions can allow 4D flat geometry to coexist with nonzero brane tension was explored\textsuperscript{6} and rejected\textsuperscript{25} within a 5D context. Given the similarity in their motivations it is natural to wonder if the 6D proposal can be killed in a similar way.
The objection in the 5D case argued that hidden fine-tunings were involved, because the presence of a brane with positive tension, $T_1$, necessarily warps the transverse dimensions and forces the bulk geometry to have a singularity elsewhere in the bulk. This singularity is naturally interpreted as the presence of a second brane, and on general grounds this second brane is found to have a negative tension, $T_2$, with 4D flatness requiring a cancellation between $T_1$ and $T_2$. How could this cancellation possibly survive integrating out heavy physics on only one brane?

The analogue of this 5D argument arises for 4D-flat compactifications of 6D supergravity on a torus. In this case using $R_2 = \chi = 0$ in eq. (2) implies $\sum_b T_b = 0$, which shows that all such compactifications require cancellation amongst brane tensions. But crucially eq. (2) does not rely on 4D flatness. Rather, being topological it must continue to hold under any continuous perturbation, such as the 'integrating out' of short-wavelength physics. If, in particular, the tension is adjusted on only one brane then eq. (2) remains true, and implies the bulk geometry must necessarily curve in response. A real calculation is required to see whether the observed 4 dimensions also curve.

What about Weinberg's Theorem?

Weinberg has a general no-go theorem against approaches to the CC problem (like SLED) which rely on spontaneously broken classical scale invariance. In a nutshell, this states that scale invariance by itself cannot protect the CC from quantum corrections, even in the absence of scale anomalies. As discussed in more detail elsewhere, as applied to SLED Weinberg's argument correctly implies that the scale-invariant classical flat direction of 6D supergravity must be lifted by quantum corrections. It does not in itself say how big these corrections must be. In SLED this lifting is partially protected by the unusually small bulk-supersymmetry breaking scale, $\Delta m^2 \sim 1/r^2$, and so must vanish as $r \to \infty$. The key work in progress for SLED is showing that this suppression is of order $1/r^4$ (in the Jordan frame) rather than merely being of order $M^4$ or $M^2/r^2$. (Notice that although $M^2/r^2$ is too large to be consistent with the observed Dark Energy, it is parameterically small compared with the normal size, $O(M^4)$, usually obtained within models.)

What is the 4D relaxation mechanism for Self-Tuning?

This objection argues against the possibility of 'self-tuning', i.e. of there being any system for which perturbations of initially 4D-flat solutions lead to new static 4D flat solutions. It arises in the following two different forms:

Volume Stabilization and Self-Tuning: One form of the argument argues that eventually it is necessary to stabilize the extra dimensions, and so develop a minimum for some sufficiently large value of $r \sim r_0$, with an effective 4D potential satisfying $V(r_0) \lesssim 1/r_0^4$. Once this is done, self-tuning would require the potential to adjust dynamically to any changes of microscopic parameters (such as a phase transition on one of the branes) in such a way as to obtain a new stabilized minimum at $r \sim r_0'$, with $r_0' \sim r_0$ and $V(r_0') \lesssim 1/r_0'^4$. But there is no known way to do this within a 4D effective description.

While this seems to be a true statement, SLED is not a self-tuning system. At the classical level this can be seen because of its flat direction: any classical perturbation stimulates a roll along the flat direction and so is not attracted towards a new static solution. Furthermore, this property can survive the lifting of the flat direction by quantum corrections, because the resulting $1/r^4$ potential typically does not stabilize the volume, and also favours a runaway along the would-be flat direction. At the 4D level the low-energy dynamics is described by a

\footnote{We thank N. Arkani-Hamed, G. Dvali and J. Polchinski for arguments along these lines.}
scalar-tensor theory, with the $1/r^4$ potential naturally giving a scalar mass of order the Hubble scale: $m \sim H \sim \mu^2/M_p$.

What is remarkable is that a light scalar with a potential of the form obtained by dimensional reduction can potentially provide a description of what is seen in cosmology. To see how this might work, imagine that the classical flat direction along which $r$ can change (parameterizing the classical scale invariance of the 6D equations) is lifted by quantum effects that are of order $1/r^4$ (which of course is the hard part, see above). Such a potential drives a runaway out to $r \to \infty$, without stabilizing at any fixed $r$. However, the full quantum contributions to the potential generically also include logarithms: $V(r) \sim (1/r^4)[a + b \log r + \cdots]$. It happens that this kind of potential can describe a successful quintessence-type cosmology, with potential domination occurring when $\log r \sim a/b$. Of course the success of this cosmology requires finding a compactification for which $a/b \sim 60$ (which is not yet done); that the universe to start off with somewhat special initial conditions (which we expect in any case from the 6D point of view); and that the light scalar which results is not ruled out by tests of gravity (also possible – but not generic – given the log $r$ corrections to its matter couplings). But these are all issues which are likely to be easier to solve than is the original cosmological constant problem.\footnote{An alternative possibility has the coefficient, $A(r)$, of the kinetic term, $A(r)(\partial r)^2$, vanishing for $r \sim r_0$.}

Adiabatic Version: An alternative version of this argument starts from the observation that there is nothing in a low-energy effective theory which precludes having a CC which is larger the cutoff, provided that it is turned on in a sufficiently adiabatic way. As applied to extra dimensional models, this appears to mean that a CC larger than $1/r^4$ could make sense in the effective 4D theory designed to describe physics at scales $E \ll 1/r$. This argues that extra dimensions cannot be key to the argument, since it should be possible to understand purely in four dimensions why one cannot add a large CC compared with the present-day Kaluza-Klein (KK) scale.

The loophole in this argument lies in its ignoring of the scale invariance of the higher-dimensional models, which preclude having a strictly constant term in the low-energy potential. Rather, any large energy density must really arise in the low-energy theory as a potential for $r$, of the form $M^{4-n}/r^n$. However, for canonically normalized fields, $\omega \sim M_p \log(Mr)$, this becomes $A e^{\lambda \omega/M_p}$, for some $\lambda$, with $A \sim M^4$ large. However, besides providing a large energy density, such a term also acts as a force pushing $\omega$, implying in particular $\dot{\omega}^2 \sim M^4$. Since this implies $\dot{\omega}$ is greater than the KK scale, it is necessarily non-adiabatic and so not describable purely within the 4D theory.

4 Summary

Applying extra dimensions to the cosmological constant problem is clearly a work in progress. However the stakes are high and include the validity of naturalness as a fundamental theoretical criterion. On the one hand extra dimensions are attractive, inasmuch as they break the basic link between vacuum energy and 4D curvature which is at the root of the CC problem. On the other hand, it is not yet known whether loop corrections can be as small as a truly natural solution to the CC problem would require (although this should be known very soon).

Moreover, even if extra dimensions provide stability under renormalization, they inevitably appear to involve special choices of initial conditions if they are to describe our present-day universe. Whether this is a reasonable trade-off, or involves throwing out the baby with the bathwater, can only be decided by examining extra dimensional solutions in more detail.
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PIERRE AUGER OBSERVATORY STATUS AND RESULTS

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The Pierre Auger Observatory, a hybrid detector for the study of ultra-high energy cosmic rays (UHECRs), is now approaching completion. After describing Auger present status and performance, with an emphasis on the advantages provided by the combination of two different detection techniques, this contribution presents a brief panorama of the first scientific results achieved and of their impact on our knowledge of the UHECRs' origin and composition.

1 Introduction

Despite the important progress achieved in cosmic ray physics during the last decades, fundamental questions about the nature and origin of the ultra-high energy cosmic rays (UHECR) are still unanswered. Contradictory results\(^1\,^2\) have been reported about the presence of the expected GZK cutoff in the cosmic ray spectrum at energies around \(5 \times 10^{19}\) eV; and the identification of possible acceleration sites still awaits the observation of an unambiguous correlation of UHECR with astrophysical objects (see\(^3\) for recent reviews on these issues). The Pierre Auger Observatory\(^4\) is expected to shed some light on these longstanding questions. Its hybrid design, combining a surface detector (SD) and a fluorescence detector (FD), makes it sensitive to different - and complementary - observables of the extensive air showers (EAS) related to the primary UHECR properties.

With more than 75\% of the SD stations deployed and all four fluorescence telescopes operational at the time of writing, the Auger Southern Site (located in the province of Mendoza, Argentina) is now nearing completion and has been accumulating high-quality data at a regularly increasing pace for the past three years. After a brief description of the detector and its current performance in Sec. 2, a review of the significant physics results already produced by Auger concerning the UHECR energy spectrum (Sec. 3), arrival directions (Sec. 4) and composition (Sec. 5) will be presented.
2 Status and description of the observatory and its dataset

The SD is a triangular array of 1600 water tanks distant 1.5 km from each other, which sample the shower content at ground. The Cherenkov light emitted by the particles entering the tank is detected by three photomultipliers and the corresponding signals are digitized at 40 MHz by Flash Analog-Digital Converters (FADC). Two local triggers are used: a simple "Threshold" (Th) one, and a "Time-over-Threshold" (ToT) one which requires a lower but more extended signal (at least 12 FADC bins) and is more sensitive to the electromagnetic (EM) component of the shower. A global trigger (T3) then asks for a relatively compact configuration of local triggers compatible in time with the arrival of a shower front. Finally, offline criteria are applied to reject accidental triggers ("physics trigger", T4) and to ensure the reconstructibility of the events ("quality trigger",T5)\(^5\). The SD is constantly active and provides the bulk of data required for high-statistics analysis. Its detection efficiency is 100% above 10\(^{18.5}\) eV at zenith angles below 60°. The angular accuracy on the arrival direction is determined on the basis of an empirical model for the time measurement uncertainties\(^6\); it depends on the number of hit stations but is always better than 2° for events at \(\theta \leq 60°\).

The SD array is overlooked by four FD sites that measure the ultraviolet light produced when charged particles in the air shower excite nitrogen molecules in the atmosphere. Each site features six Schmidt telescopes that cover a field of view of about 30° × 30° each. The signal is collected on a 440 pixels camera and digitalized at a 100 MHz sampling rate. The fluorescence light emitted by the shower is roughly proportional to the energy dissipated in the atmosphere. The fluorescence telescopes can be used only during dark, moonless nights, which reduce their duty cycle to about 14%. The timing and position of the triggering pixels allow to reconstruct the shower-detector plane with an accuracy of about 0.3°, but the uncertainty on the orientation of the shower axis within that plane is much larger.

Most events seen by the FD also trigger at least one SD station, and the additional timing information allows to significantly improve the accuracy both on the reconstructed arrival direction (\(\sim 0.5°\)) and the position of the core (\(\sim 50m\)). These hybrid events amount to about 10% of the total data sample; they allow to fully exploit the detector capabilities and therefore have an important impact on many analysis performed on Auger data. High-quality events, which independently trigger the FD and the SD and can be successfully reconstructed by both detectors, are tagged as golden hybrids. These events allow the simultaneous measurement and cross-calibration of different shower parameters related to the energy and nature of the UHECRs. Such a strategy allows to extract physical information about their spectrum and composition while minimizing the dependance in model assumptions, as will be illustrated in the next sections.

3 The spectrum of UHECR

3.1 General strategy

The key ingredients for the determination of the UHECR spectrum are the accurate determination of the primary energy, which is best achieved with the fluorescence technique, and a large and easily calculable exposure, which is provided by the SD.

The signals in the triggered stations are used to reconstruct both the shower core position and its lateral profile at ground. The parameter \(S(1000)\), i.e. the signal that would be produced in a tank located at 1000 m from the shower core, is measured with an accuracy better than 12% and can be used as an estimator of the size (and thus energy) of the shower\(^8\). For a given energy, its value depends on the zenith angle of the shower as a consequence of geometrical effects and of the attenuation of the shower in the atmosphere. The "constant intensity cut" method\(^7\) is used to extract the shape of this attenuation curve, \(CIC(\theta)\), from the data assuming
an isotropic flux of UHECRs. The $S(1000)$ is then converted into a reference value taken at the mean of the zenith angle interval, $S_{380} \equiv S(1000)/CIC(\theta)$.

The relation of $S_{380}$ (or $S(1000)$) to the primary energy however significantly depends on the assumptions on the primary composition and on hadronic models which drive the development of the shower. This drawback is circumvented by using the golden hybrid events to calibrate $S_{380}$ on the energy obtained with the FD. The information on the shower longitudinal profile provided by the FD indeed allows an independent, nearly-calorimetric measurement of the energy of the shower, $E_{FD}$. Dependence in composition and hadronic models only affects the determination of the invisible component, \textit{i.e.} muons and neutrinos, which contributes only 4% of the total uncertainty in the FD energy. More significant sources of systematics are the uncertainty on the fluorescence yield and its dependence in the atmospheric conditions, the absolute calibration of the FD and the energy reconstruction method. Current estimations\textsuperscript{9} of the overall systematics in $E_{FD}$ give about 22%, while the statistical uncertainty in the derived energy is smaller than 10%.

The dataset now used to build the spectrum includes all SD T5 events recorded between January 1st, 2004 and February 28th, 2007, with reconstructed $\theta \leq 60^\circ$ and energy $E_{FD} \geq 3$ EeV, which ensures full detection efficiency of the SD and allows a geometrical computation of the corresponding aperture. After removal of periods of failure in data acquisition, the corresponding integrated exposure amounts to about 5165 km$^2$ sr yr, which is more than three times the one obtained by the AGASA experiment \textsuperscript{1}. The corresponding spectrum is shown in Fig. 1 together with its statistical uncertainty. Although the statistics is still limited, the hypothesis of a continuation of the UHECR spectrum in the form of a pure power law beyond $10^{18.6}$ eV can now be rejected at a 6$\sigma$ confidence level, as discussed in \textsuperscript{11}. Efforts to reduce the systematics in the energy estimation are in progress as well. In particular, recent and ongoing measurements\textsuperscript{12} of the fluorescence yield at a precision level of 5% are expected to significantly improve the accuracy in the reconstructed $E_{FD}$.

3.2 Spectrum from inclined events

The use of Cherenkov water tanks as surface detectors allows the Auger to detect showers with zenith angles up to $90^\circ$ (and even beyond). The range of inclined showers, $60^\circ \leq \theta \leq 90^\circ$,
contributes half the total solid angle of the detector and about 25% of its geometrical acceptance, thereby significantly increasing the field of view of the detector and the SD statistics.

Such showers are characterized by a dominance of the muonic component at ground and by a very elongated and asymmetrical footprint which can exhibit a lobular structure due to the bending action of the geomagnetic field. The energetic (10 – 1000 GeV) muons reach the detector in a thin front with small curvature, which produces short and peaked FADC pulses. Dedicated selection procedures and reconstruction methods, based on the use of density maps of the number of muons at ground, have been developed to analyze such events; more detail can be found in \cite{14}.

The strategy for building a spectrum is the same as in Sec. 3.1. Once the arrival direction is reconstructed, the pattern of signals is fitted to muon density maps obtained from simulated proton showers at $10^{19}$ eV in order to determine the core position and an overall normalisation factor, $N_{19}$, which acts as an energy estimator and can be calibrated on the FD energy using inclined hybrid events. No constant intensity cut is needed because the muon maps already account for the shower attenuation and geometrical effects. Inclined events with $60^\circ \leq \theta \leq 80^\circ$ and $E_{FD} \geq 6.3$ EeV ($N_{19} \geq 1$), where the SD detection efficiency is expected to be 100%, have been used to build this independent spectrum \cite{15}. The corresponding integrated exposure amounts to 1510 km$^2$sr yr, about a quarter of that of "vertical" ($\theta \leq 60^\circ$) events.

3.3 Spectrum from hybrid events

Although their statistics is much smaller, hybrid events alone allow a spectrum determination below the energy threshold of the SD \cite{13}. To guarantee the quality of the reconstruction, only events with a reconstructed $\theta \leq 60^\circ$ and satisfying extra requirements on the observed profile were selected. In particular, the contamination by Cherenkov light may not exceed 50% and the reconstructed depth of the shower maximum, $X_{max}$, must be observed and lie within a fiducial volume (which depends on the energy) in order to avoid biases due to the limited field of view.

The hybrid exposure is estimated on basis of a detailed Monte Carlo simulation which accounts for the growth of both FD and SD during the data taking period, as well as for seasonal and instrumental effects. The hybrid trigger efficiency reaches 100% at $10^{18}$ eV, independently of the nature of the primary (proton/iron). The main sources of uncertainty again lie in the determination of the energy (and its impact on the event selection and aperture calculation), the knowledge of atmospheric conditions and the estimation of the detector uptime; see \cite{13} for a more detailed discussion.

The three spectra (multiplied by $E^3$) are compared in Fig. 2. The spectrum from SD vertical events is the most accurate and statistically significant. The hybrid spectrum extends to lower energies and encompasses the "ankle", which appears as a spectral break at $\sim 10^{18.5}$ eV. A detailed assessment of the sources of systematics remains to be done for inclined events, but all three spectra are in reasonable agreement within current estimated uncertainties.

4 The arrival direction of UHECR: anisotropy searches

Anisotropies in the flux of UHE cosmic rays may appear in different energy ranges and angular scales, depending on the nature, distance and extension of the source(s). Cosmic rays around an EeV are thought to be of galactic origin, and the region of the Galactic Center (GC) and the Galactic Plane (GP) are key targets for anisotropy searches performed with Auger data in that energy range. At higher energies one rather expects UHE cosmic rays to come from extragalactic sources; a search for directional excesses of cosmic rays could then reveal a correlation with astrophysical objects or even exotic sources.
The anisotropy studies performed by Auger are based on both SD T5 and hybrid events with $\theta < 60^\circ$ and the energy assigned via the cross-calibration procedure described in Sec. 3.1.

### 4.1 Anisotropy studies around the Galactic Center and the Galactic Plane

In the past, two cosmic ray experiments, AGASA and SUGAR, have claimed significant excesses in the flux of UHECR in the region of the GC\textsuperscript{16,17}. Recent TeV $\gamma$-ray observations by HESS\textsuperscript{18,10} have provided additional hints towards the presence of powerful CR accelerators in the Galaxy. In that context, several models that predict a detectable flux of neutrons from the GC in the EeV range (when the neutron decay length is about the distance from the GC to the Earth) have also been proposed\textsuperscript{20}.

With the GC well in the field of view and a much better angular resolution than previous CR experiments, the Pierre Auger Observatory is well suited to look for UHECR anisotropies coming from that region. Such a search was performed on the bulk of SD data in different energy ranges and with different sizes of the angular filtering in order to match the resolution of previous experiments. Using a data sample much larger than the AGASA and SUGAR ones (79265 SD events and 3934 hybrid events with $\theta < 60^\circ, 10^{17.0}$ eV $< E < 10^{18.5}$ eV, corresponding to the period from January 2004 to March 2006), Auger did not confirm any of the anisotropy claims\textsuperscript{21}.

The same data were also used to search for a point source in the direction of the GC itself at the scale of Auger's own angular resolution\textsuperscript{21}. Applying a 1.5° Gaussian filter to account for the pointing accuracy of the SD, no excess of events was observed. Assuming both the source and the bulk CR at those energies have a spectrum index of 3.3 and that the emitted CRs are protons, an 95% C.L. upper limit of $\Phi_x^0 \leq \xi$ 0.13 km$^{-2}$ yr$^{-1}$ (where $\xi$ parameterizes the uncertainties on the flux normalization) is set on the source flux\textsuperscript{21}.

A recent update of this analysis with a better angular accuracy and a significantly larger dataset allowing to split the energy range in 0.1 EeV $\leq E \leq 1$ EeV and 1 EeV $\leq E \leq 10$ EeV have confirmed all negative anisotropy results and improved the bound on $\Phi_x^0$ to $\xi$ 0.018 km$^{-2}$ yr$^{-1}$. Such a limit already excludes most of the models of neutron production at the GC\textsuperscript{22}.

Finally, several methods have been set up to search for large-scale anisotropies in the distribution of UHECR at energies around the EeV (and above); such angular patterns would hint towards a galactic origin of the UHECR just below the ankle. With the current data set, the right ascension distribution is found to be compatible with an isotropic sky\textsuperscript{23}. Searches for bidimensional patterns, such as a possible dipole, are also ongoing.

### 4.2 Searches for localized excesses and correlations with astrophysical objects

Blind searches using Auger data have been performed looking for small- and intermediate-scale excesses in the sky that would reveal the presence of point-like sources. The statistical significance of such an excess is estimated by calculating the two-point angular correlation function, which counts the number of pairs of events with energy larger than a given threshold $E_{th}$ separated by less than an angle $\theta$. Recent studies on SD T5 data with $E > 10$ EeV, scanning a large range of ($\theta, E_{th}$), shows no really significant signal of anisotropy, although some hints of clustering exist at very high energies and intermediate angular scale\textsuperscript{24}.

Events with energies above 10 EeV have also been used to test a possible correlation with subsets of BL Lacs, in relation with previous (and sometimes contradictory) claims and results based on data from AGASA, Yakutsk and HiRes experiments\textsuperscript{25}. With 6 times more events than the other existing data samples, the analyzed Auger data is still compatible with isotropy and does not support any of the previously reported signals of clustering.

\textsuperscript{21}This bound could however be about 30% higher if the CR composition at EeV were heavy.
Anisotropy searches based on Auger list of prescribed targets with definite angular and energy windows as released in\textsuperscript{26} has also given negative results. As more data is streaming in, the catalogue of candidate targets that will be studied is expected to increase in the future.

5 The nature of UHECR: composition studies

Thanks to its hybrid design, Auger can in principle measure an extended set of parameters sensitive to the primary UHECR nature and mass. While the discrimination between different types of nuclei is complicated by the uncertainties in the hadronic models, several methods have already been proposed for the identification of photon and neutrinos. The detection of such particles in the UHE cosmic radiation would probe many exotic models of UHECR production and help locate candidate sources as they travel undeflected by the ambient magnetic fields.

5.1 Upper limit on the flux of UHE photons

Unlike protons and nuclei, the development of photon showers is driven by electromagnetic (EM) interactions and does not suffer much from the uncertainties in hadronic models. Their development is also delayed due to the small multiplicity in EM interactions and to the LPM effect\textsuperscript{27}, which reduces the bremsstrahlung and pair production cross-sections above 10 EeV.

One of the methods set up by Auger to identify photon primaries in the flux of UHECR is based on the direct observation of the longitudinal profile of the shower by the FD; the discriminating variable is the atmospheric depth of the shower maximum, $X_{\text{max}}$ (the estimated average difference in $X_{\text{max}}$ between photons and hadrons is about 200 g/ cm$^2$). The data set used for this analysis are hybrid events with a reconstructed energy $E > 10^{19}$ eV registered between January 2004 and February 2006. A series of cuts were applied to guarantee the quality of the hybrid geometry and of the fit to the shower longitudinal profile (which takes into account the local atmospheric conditions), and to minimize the bias against photons introduced in the detector acceptance by requiring the $X_{\text{max}}$ to be inside the field of view. For all (29) events passing the cuts, the observed $X_{\text{max}}$ is well below the average value expected from the simulation of 100 photons showers in the same conditions. Taking systematic uncertainties on the $X_{\text{max}}$ determination and the photon shower simulations into account, this analysis, described in\textsuperscript{28}, allowed to put a 95% C.L. upper limit on the photon fraction of 16% above 10 EeV; it has been recently updated to 13% using a more extended data sample\textsuperscript{29}, as shown in Fig. 3.

Another analysis relying on the SD measurements has also been developed; the key observables are here the signal risetime at 1000 m (i.e. the time it takes for the signal to rise from 10% to 50%) and the radius of curvature of the shower front. Particles from showers with a larger $X_{\text{max}}$ (and thus a later development) are indeed expected to reach the ground in a thicker and more curved front. A principle component analysis combining both observables was used to search Auger data for photons; no candidate was found and the corresponding upper limit on the photon fraction is 2.0%, 5.1% and 31% at 10, 20 and 40 EeV respectively.

As shown in Fig. 3, the stringent limits put by Auger results on the UHE photon fraction now disfavour many of the top-down models proposed in connection with the AGASA spectrum.

5.2 Upper limit on the flux of UHE neutrinos

Due to their small interaction cross-section, neutrinos can penetrate large amounts of matter and generate showers at any atmospheric depth, unlike protons or photons. Young and deep neutrino-induced showers can thus be efficiently identified in the range of inclined showers, $\theta \geq 60^\circ$, by requiring the presence of a significant EM component.

Upward-going tau neutrinos that graze the Earth just below the horizon could also be detected as they are likely to interact in the crust and produce a tau lepton which may emerge and
Figure 3: Auger upper limits on the UHE photon fraction from the hybrid analysis (labeled FD) and from SD analysis (black arrows), together with some predictions from top-down models and the bounds put by previous experiments (from [33]).

Figure 4: Auger upper limit on an $E^{-2}$ diffuse flux of UHE $\nu_\tau$ (with the worst estimation of systematics), together with some predictions for GZK neutrino fluxes and the bounds put by other cosmic neutrino experiments (from [32]).

initiate an observable shower, provided it does not decay too far from the ground. This channel had been pointed out as likely to increase the detection potential of Auger for neutrinos in the EeV range [30], and it has been extensively studied in a recent analysis [31, 32].

The emerging $\tau$ flux corresponding to a given incident $\nu_\tau$ flux has been computed in the relevant angular window using both Monte Carlo and semi-analytical methods accounting for all $\nu_\tau \leftrightarrow \tau$ conversion processes, as well as for the $\tau$ energy losses. The atmospheric shower produced by the decay of the emerging $\tau$ is then simulated and tested for detection in the SD.

A specific selection procedure has been set up to identify those $\tau$-induced, nearly horizontal showers that develop close to the detector. The signal shape must be compatible with the presence of a significant EM component (in practice a ToT-type trigger is required) and a large (>1.4) area-to-peak ratio to reject triggers produced by consecutive muons. The footprint of those stations is then required to assume an elongated shape and the timings to be compatible with a shower front traveling nearly horizontally at the speed of light. The efficiency of identification of a $\tau$-induced shower depends on $E_\tau$ and on $h_{\text{c}}$, the altitude of the shower center (defined at a nominal distance of 10 km from the $\tau$ decay point along the shower axis), but also on the relative position of the footprint in the array. To compute the detector acceptance for $\nu_\tau$, a double Monte Carlo integration accounting for the evolution of the array with time is performed.

SD data from January 2004 till December 2006 were searched for candidate grazing $\nu_\tau$'s, but no single event passed the selection criteria, which allows to derive an upper limit for any injected flux of UHE $\nu_\tau$ with a given shape. Assuming an $E^{-2}$ incident spectrum of diffuse $\nu_\tau$, a 90% C.L. bound $E_\nu^2dN_\nu/dE_\nu < 1.5_{-0.8}^{+1.5} \times 10^{-7}$ GeV cm$^{-2}$sr$^{-1}$s$^{-1}$ was derived in the energy range [2 $10^{17}$ − 5 $10^{19}$] eV. The sources of systematic uncertainties have been carefully addressed [32]; they are globally responsible for a factor of ~3 uncertainty on the acceptance, which propagates to the final flux limit. Among them, physical quantities that have not been measured at those energies, such as the $\nu$ cross-sections, the $\tau$ polarization and energy losses, contribute resp. ~15%, ~30% and ~40%. The Monte Carlo simulations of the shower and the detector response add an extra ~25% uncertainty, and the effect of neglecting the actual topology of the Auger site another ~18%. As shown in Fig.4, Auger limit is nevertheless the best to date in the energy range where GZK neutrino fluxes (produced by the interaction of the observed UHECR with the cosmic microwave background) are likely to peak. To improve that limit by an order of magnitude or so will however require the accumulation of several more years of data.
6 Conclusions

The past three years have witnessed a phase of major development of the Southern Auger Observatory on the field, accompanied by a significant increase of the dataset. A lot of progress has been made in the understanding the detector, which resulted in a better control on the systematic uncertainties an in the development of reliable and robust analysis methods which allowed the release of first scientific results concerning the UHECR spectrum and angular distribution. If a continuation of the spectrum above $10^{20}$ eV seems unlikely, a much larger data sample is still needed to determine the exact shape of the spectrum. Auger also sees a remarkably isotropic sky, except maybe at high energies where more data are necessary for a detailed study of clusters and correlations, whatever the scale. Finally, Auger has already put competitive limits on the fluxes of UHE photons and neutrinos, thereby demonstrating its capabilities to work as a multi-messenger detector.

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References

ANISOTROPY OF DARK MATTER ANNIHILATION IN THE GALAXY

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Dark matter (DM) annihilation in the Galactic halo is strongly enhanced (boosted) with respect to a diffuse DM annihilation by the presence of small-scale DM clumps. The distribution of clumps in the Galactic halo is described in the framework of standard cosmology and hierarchical structure formation by taking into account a tidal destruction of clumps by stars. A tidal destruction of clumps in the Galactic disk results in an anisotropy in clump distribution. A corresponding annihilation of dark matter particles in small-scale clumps produces the anisotropic gamma-ray signal with respect to the Galactic disk. This anisotropy is rather small, $\sim 9\%$, and superimposed on that due to off-centering position of the Sun in the Galaxy. The anisotropy of annihilation signal with respect to the Galactic disk provides a possibility to discriminate DM annihilation from the diffuse gamma-ray backgrounds of other origin.

1 Introduction

A primordial power-law spectrum of density fluctuations in the Dark Matter (DM) ranges from the largest scales above the scales of superclusters of galaxies to the smallest sub-stellar scales according to prediction of inflation models. This permits to predict the properties of smallest DM structures from the known CMB fluctuations at large scales. Substructures of DM in the galactic haloes with a rather large mass, $\geq 10^7 M_\odot$, were extensively discussed in early works, see for example\textsuperscript{1}. The nonlinear dynamics and mechanism of hierarchical clustering of these large DM clumps were analyzed in both analytical calculations\textsuperscript{2,3,4} and numerical simulations\textsuperscript{5,6,7}. At sub-stellar mass-scales of DM fluctuations, a principal new phenomenon arise — the cutoff of mass spectrum due to collisional and collisionless (free streaming) damping processes of DM particles in the forming clumps. The resulting smallest mass of DM clumps is determined by the properties of DM particles, in particular, by their elastic scattering. See e. g.\textsuperscript{8} and references therein for detailed calculations of this cutoff. Additionally the cutoff of mass spectrum is influenced by
the acoustic absorption\(^9\) at the time of kinetic decoupling of DM particles\(^10\) and also by the horizon-scale perturbation modes\(^11\). The kinetic equations for DM phase space density were solved\(^12\) for the case of perturbed cosmological background by taking into account the acoustic absorption, horizon-scale modes and gravitational perturbations. A corresponding value of the smallest clump mass for neutralino DM is of the order of the Moon or Earth mass. The formation of small-scale DM clumps with a mass larger than the Earth mass, \(M_{\min} \sim 10^{-6} M_\odot\), have been explored in numerical simulations\(^13,14\). A resulting differential number density of small-scale clumps, \(n(M) dM \propto dM/M^2\), turns out very close to that obtained in the numerical simulations of large-scale clumps with mass \(M \geq 10^6 M_\odot\). The other important result obtained in numerical simulations\(^13\) is determination of the internal density profile in the isolated clump of minimal mass. The resulting density profile is approximately a power-law, \(\propto r^{-\beta}\), with \(\beta = 1.5 - 2.0\), which is in a good agreement with theoretically predicted value \(\beta = 1.7 - 1.8\) according to\(^2\).

The number density of small-scale DM clumps existing nowadays in the Galactic halo is determined by their tidal destruction during hierarchical structure formation\(^15\) and also by tidal interactions with stars in the Galaxy\(^16,17,18,19,20,21\). Annihilation of DM particles in small-scale clumps\(^14,22,23,24,25,26,27,28,29\) enhances the total DM annihilation signals in our Galaxy and thus boosts a chance for indirect detection of DM.

The usual assumption in calculations of DM annihilation is a spherical symmetry of the Galactic halo. In this case an anisotropy of annihilation gamma-radiation is only due to off-center position of the Sun in the Galaxy. Nevertheless, a principal significance of the halo nonsphericity for the observed annihilation signal was demonstrated in\(^30\). According to observations, the axes of the Galactic halo ellipsoid differ most probably no more than \(10 - 20\%\), but even a much more larger difference of axes, up to a factor 2, can not be excluded\(^31,32\). This leads to more than an order of magnitude uncertainty in the predicted annihilation flux from the Galactic anti-center direction\(^30\). It must be noted also the “intrinsic” annihilation anisotropy caused by the small-scale DM clustering itself. A corresponding angular power spectrum of annihilation signal at small scales is connected with a power spectrum of DM clumping\(^33\). In principle, the DM clumps may be seen as point sources at the gamma-sky\(^30\). Another minor source of annihilation anisotropy is a dipole anisotropy due to proper motion of the Sun in the Galaxy\(^34\).

In\(^18\) the anisotropy with respect to the Galactic disk was discussed basing on the numerical calculations of the destruction of DM clumps by stars in the disk and taking into account the influence of gravitational potential of the disk on the clump orbits. It was also shown\(^15,19\) that (i) small-scale DM clumps dominate in the generation of annihilation signal and (ii) the Galactic stellar disk provides the main contribution to the tidal destruction of clumps at \(r > 3 \text{ kpc}\), i.e. outside the central bulge region. A process of clump destruction in the halo is anisotropic in general (e.g. it depends on the inclination of clump orbit with respect to the disk plane). Respectively, the DM annihilation in the halo is also anisotropic. In this work we estimate the value of this anisotropy. It must be stressed that with a present state of art it is impossible to separate this source of anisotropy from that produced by the halo nonsphericity. More detailed investigation is required to constrain the shape of the halo and to search the distinctive features of annihilation anisotropy due to non-spherical halo clumpiness. The detectors at the GLAST satellite will be sensitive to anisotropy up to \(0.1\%\) level\(^34\). This will provide a hope to discriminate the anisotropic DM annihilation signal from the diffuse gamma-ray backgrounds.

2 Anisotropic destruction of clumps by disk

Crossing the Galactic disk, a DM clump can be tidally destructed by the collective gravitational field of stars in the disk. This phenomenon is similar to the destruction of globular clusters by the “tidal shocking” in the Galactic disk\(^35\). The corresponding energy gain per unit mass of a
clump at one disk crossing\(^{35}\) is
\[
\Delta E = \frac{2g_m^2 (\Delta z)^2}{v_{z,c}^2},
\]
where \(g_m\) is the maximum gravitational acceleration of the clump moving through the disk, \(\Delta z\) is a vertical (perpendicular to the disk plane) distance of a DM particle from the clump center, \(v_{z,c}\) is a vertical component of velocity at disk crossing. The dependence of \(v_{z,c}\) on the inclination of orbit relative to the disk plane is the origin of the discussed anisotropy in the clump destruction, and, as a result, the origin of the anisotropy in annihilation signal.

The surface mass of the Galactic disk can be approximated as
\[
\sigma_s(r) = \frac{M_\odot}{2\pi r_0^2} e^{-r/r_0},
\]
with \(M_\odot = 8 \times 10^{10} M_\odot\) and \(r_0 = 4.5\) kpc, and therefore
\[
g_m(r) = 2\pi G \sigma_s(r).
\]
We use the power-law parametrization\(^{2,3,4,13}\) of the internal density of a clump
\[
\rho_{\text{int}}(r) = \frac{3 - \beta}{3} \rho \left( \frac{r}{R} \right)^{-\beta},
\]
where \(\rho\) and \(R\) are the mean internal density and a radius of clump, respectively, \(\beta = 1.8\) and \(\rho_{\text{int}}(r) = 0\) at \(r > R\). The total (kinetic plus potential) internal energy of a clump for density profile (4) is given by
\[
|E| = \frac{3 - \beta}{2(5 - 2\beta)} \frac{GM^2}{R},
\]
where \(M\) is the mass of the clump. Integrating (1) over a clump volume and using the density profile (4), one obtains an energy gain for the whole clump as
\[
\frac{\Delta E}{|E|} = \frac{(5 - 2\beta)}{\pi(5 - \beta)} \frac{g_m^2}{G \rho v_{z,c}^2}.
\]
We will use the following criterium for a tidal destruction of clump: a clump is destructed if a total energy gain \(\sum \Delta E_i\) after several disk crossings exceeds the initial internal energy \(|E|\) of a clump.

Let us consider now some particular orbit of a clump in the halo with an “inclination” angle \(\gamma\) between the normal vectors of the disk plane and orbit plane. The orbit angular velocity at a distance \(r\) from the Galactic center is \(d\phi/dt = J/(m_r r_c^3)\), where \(J\) is an orbital angular momentum of a clump. A vertical velocity of a clump crossing the disk is
\[
v_{z,c} = \frac{J}{m_r r_c} \sin \gamma,
\]
where \(r_c\) is a distance of crossing point from the Galaxy center. There are two crossing points (with different values of \(r_c\)) during the one orbital period. The momentum approximation used here for calculations of the tidal heating is violated at small inclination angles, \(\gamma \ll 1\). Nevertheless the resultant anisotropy is a cumulative quantity. It results from an integration over all clump orbits, and orbits with \(\gamma \ll 1\) provide only small input into the anisotropy value.

A tidal heating and final destruction of clumps by the gravitational field of the Galactic disk depends on the inclination angle \(\gamma\) of a clump orbit to the disk according to (1). This is a cause of the anisotropic clump number density decreasing during the lifetime of the Galaxy. The numerically calculated survival probability of DM clumps in the halo\(^{36}\) is shown in the Fig. 1. The annihilation anisotropy is artificially enhanced in the Fig. 1 for better visualization for three chosen radial distances from the Galactic center, \(r = 3, 8.5\) and 20 kpc respectively by using the different multiplication factors.
3 Annihilation anisotropy

For the diffuse distribution of DM in the halo, the annihilation signal (e.g., gamma-ray or neutrino flux per unit solid angle) is proportional to

\[
I_H = \int_0^{\tau_{\text{max}}(\zeta)} \rho_H^2(\xi) \, dx,
\]

where \( x = r/L \) and integration over \( r \) goes along the line of sight, \( \xi(\zeta, r) = (r^2 + r_\odot^2 - 2rr_\odot \cos \zeta)^{1/2} \) is the distance to the Galactic center, \( \tau_{\text{max}}(\zeta) = (R_H^2 - r_\odot^2 \sin^2 \zeta)^{1/2} + r_\odot \cos \zeta \) is the distance to the external halo border, \( \zeta \) is an angle between the line of observation and the direction to the Galactic center, \( R_H \) is a virial radius of the Galactic halo, \( r_\odot = 8.5 \) kpc is the distance between the Sun and Galactic center. The corresponding signal from annihilations in DM clumps is proportional to the quantity \(^{15}\)

\[
I_{cl} = \mu \Sigma \rho \int_0^{\tau_{\text{max}}(\zeta)} \rho_H(\xi)P(\xi, \alpha)P_{sp}(\xi) \, dx,
\]

where \( \mu \simeq 0.05 \) is a fraction of the DM mass in the form of clumps, \( P_{sp} \) is a survival probability of clumps due to their tidal destruction by stars in the halo and bulge from \(^{10}\). The function \( S \) depends on the clump density profile and core radius of clump\(^{15}\) and we use \( S \simeq 14.5 \). Here for simplicity we do not take into account the distribution of DM clumps over their internal densities. As a representative example we consider the Earth-mass clumps \( M = 10^{-20} M_\odot \) originated from 2\sigma density peaks in the case of power-law index of primordial spectrum of perturbations \( n_p = 1 \). The mean internal density of these clumps is \( \rho \simeq 7 \times 10^{-23} \text{ g cm}^{-3} \). The values of \( \mu \) and \( S \), as
well as the distribution of the clumps over various parameters influence the annihilation signal but only weakly affect the predicted anisotropy.

In the Fig. 2 the annihilation signal calculated according to (9) is shown for the Galactic disk plane and for the orthogonal vertical plane (passing through the Galactic center) as function of angle $\zeta$ between the observation direction and the direction to the Galactic center. For comparison in the Fig. 2 is also shown the signal from the spherically symmetric Galactic halo without the DM clumps (8). The later signal is the same in the in the Galactic disk plane and in vertical plane and therefore can be principally extracted from the observations. The difference of signals in two orthogonal planes at the same $\zeta$ can be considered as an anisotropy measure. Defined as $\delta = (I_2 - I_1)/I_1$, it has a maximum value $\delta \simeq 0.09$ at $\zeta \simeq 39^\circ$.

4 Discussions

A total anisotropy of DM annihilation signal is determined in general by the Sun off-centering in the Galaxy and by the halo nonsphericity. The small-scale DM clumps are completely destructed inside the Galactic stellar bulge region. The “gamma-rings” are predicted in other galaxies due to the absence of clumps in their centers\textsuperscript{18}. The unknown nonsphericity of the halo is a main source of anisotropy uncertainty. The value of anisotropy due to nonsphericity of the halo may be several times larger than one caused by the discussed in this paper effect of tidal destruction of DM clumps by the disk. More detailed analysis is required to separate these two sources of anisotropy. A nonsphericity (oblateness) of the halo due to the angular momentum can be easily estimated. It is natural to assume that the DM halo and disk have the same value of specific angular momentum (i.e., an angular momentum per unit mass). In this case the model of the Maclaurin spheroid for the halo gives only $\sim 0.5\%$ difference for the halo axes. Therefore, the nonsphericity of the halo due to the angular momentum produces a negligible anisotropy. The
anisotropy with respect to disk plane in the Galaxy was pointed out in\cite{13}. As it is seen from our
calculations (see Fig. 2) this anisotropy of annihilation signal from the DM clumps with respect
to the disk is rather small, $\sim 9\%$, but far exceeds the anticipated GLAST resolution, $\sim 0.1\%$.
Therefore, the discussed anisotropy may be used in future detailed gamma-ray observations for
discrimination of the annihilation signals from the DM clumps, diffuse DM in the Galactic halo
and diffuse gamma-ray backgrounds.

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Highlights from VHE Gamma Astronomy:  
Where do we stand and where do we go?

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We're witnessing the dawn of VHE Gamma Astronomy thanks to a new generation of Cherenkov Telescopes which is producing a plethora of discoveries and new measurements with a direct implication for astrophysics, fundamental physics and cosmology. The main present results and future prospects are discussed.

1 Introduction

VHE Cosmic gamma-ray observation is at present in a very special moment since a true revolution in the consolidation of Cherenkov Telescopes as astronomical instruments is taking place. After many years of slow development, Imaging Air Cherenkov Telescopes (IACTs) are now in the phase transition from being "high energy experiments" to being "telescopic installations" in the astronomical sense. This fact is motivating an exploding interest in a broad scientific community embracing astrophysics, particle physics and cosmology.

The reason for this phase transition is the big observational step occurred within the last couple of years at the quantitative level (tripling the number of detected sources) but also at the qualitative level (producing extremely high quality detections allowing unprecedented detailed studies) due to the start into operation of the Cherenkov Telescopes of the new generation.

Nevertheless, given the limited extension of this manuscript, among the broad spectrum of new results obtained from the observations in the VHE gamma ray band with the new generation of Cherenkov Telescopes we shall concentrate in discussing three highlights with the largest impact in fundamental physics and cosmology.

For that, the outline of this write-up is as follows: in section 2 we'll present the status of the indirect search for Dark Matter annihilation into VHE gamma rays discussing the impact of the detailed analyses of the Galactic center observations carried out by HESS and MAGIC. In section 3 we'll discuss the implications of the studies of the energy spectrum measured in distant Blazars by HESS and MAGIC, which allow to place an unexpectedly low upper bound on the density of the Extragalactic Background Light by means of the analysis of the gamma-gamma absorption. Finally section 4 will review the use of light curves showing fast flares of VHE gamma ray sources at cosmological distances to place constraints on the quantum structure of the gravitational vacuum.

2 Indirect searches for Dark Matter

If most of the Dark Matter is in form of Weakly Interacting Massive Particles (WIMP) as the $\Lambda$CDM scenarios, favored by most of the observations, presently suggest, a favorite candidate
for this Dark Matter is the LSP which in most supersymmetric extensions of the Standard Model is the so-called neutralino, the spin 1/2 supersymmetric partner of the neutral bosons. In that case, Dark Matter may be detected from the neutralino annihilation into pairs of VHE gamma rays from the center of our galaxy, nearby galaxies, low surface brightness dwarf spheroidal galaxies, globular clusters or “hidden” dark matter satellites.

Out of all the possible target candidates for the indirect detection of Dark Matter through its annihilation into VHE gamma rays the one from which larger flux is expected is the center of our galaxy. The reason is the very high Dark Matter density expected, which enters quadratically in the gamma ray flux prediction, and its proximity when compared to other target candidates.

The galactic center has been independently observed by HESS and MAGIC (in this case at large zenith angle, which implies larger effective area but at a higher energy threshold) providing spectrum measurements in nice agreement, which contradict the measurements previously published by the CANGAROO collaboration.

The signal observed by HESS in 2003-2004 was consistent with point-like emission from Sgr A* although it had an slight hint for extension which could be fit with a Navarro-Frenk-White Dark Matter halo profile as can be gleaned from figure 1. In addition, the signal was steady from year to minute scales. The spectrum obtained with the first data taken extended up to energies above 10 TeV and would have required invoking an unnaturally heavy neutralino to explain it.

The final spectrum after analyzing all the accumulated data can be seen in figure 2 which shows that it can be very well fitted by an unbroken power law with index 2.3 from about 150 GeV up to almost 30 TeV. This spectrum is in perfect agreement with the one obtained by the MAGIC collaboration which has observed the Galactic Center at large zenith angles (from 58 to 62 degrees) and hence, with somewhat different systematic uncertainties.

This spectrum shape and index are in agreement with the expectations from acceleration mechanisms in standard astrophysical sources and rules out most of the possible interpretations in terms of Dark Matter annihilation.

Nevertheless, a plausible explanation at that stage was that the signal from the Dark Matter annihilation in the Galactic Center region could be outshined by the VHE gamma ray emission from point-like astrophysical sources in the Galactic Center region which, from observations in many wavelengths is known to be a very busy region with many astrophysics sources and a lot of non-thermal activity.

Following this idea, HESS has been able to subtract from a deep exposure the point-like
sources (given its point-spread-function)\(^5\) and the observed remaining signal turns out to be in
good agreement with the distribution of the molecular gas traced by its CS emission as can be
observed in figure 3.

Therefore, the high quality data on the Galactic Center obtained in the last few years by
HESS and MAGIC does not show any evidence of Dark Matter annihilation signal\(^6\),\(^7\). In spite
of that, it is very difficult to extract any quantitative conclusion of that observation since there
are very large uncertainties in the predictions for the expected flux coming from:

- WIMP mass spectrum and couplings which should be known to determine the annihilation
  probabilities into the different channels. For these quantities, important accelerator and
  relic density constraints exist already but there is still a very broad parameter space
  open, which make predictions very uncertain. The start of LHC operation in the coming
  years may narrow down drastically the parameter space and allow for much more precise
  predictions.

- The cuspy region of the dark matter density profile, in the vicinity of the central super-
  masive Black Hole, which remains virtually unknown.

- The background due to astrophysical sources which may be much larger than the Dark
  Matter annihilation signal making the subtraction very uncertain

Nevertheless, other target candidates, such as Dwarf Spheroidal Satellites of our Galaxy with
high mass-to-light ratios which in comparison with the Galactic Center are expected to produce
lower fluxes and are more distant, but which may provide cleaner environments with much less
astrophysical source backgrounds, are being explored.

An important step in this search for Dark Matter annihilation signals will be the sky survey
catalog which will be produced by GLAST in the near future since its unidentified sources may
spot Dark Matter clumps and therefore be prime candidates to study in depth with Cherenkov
Telescopes in the quest for Dark Matter.
At any rate, it must be stressed that even if WIMP candidates are found in accelerator experiments it must be confirmed that they actually are constituents of the Dark Matter of our universe and for this purpose IACTs are among the most promising instruments.

3 The Cosmological Gamma Ray Horizon

As it is very well know the intergalactic vacuum is not really empty. There is a sea of photons lying around which constitute the so-called Extragalactic Background Light (EBL). For instance, one can find the well studied Cosmic Microwave Background but there are contributions from any photon energy.\(^8\)

The flux of high energy gamma rays that travel through the universe is attenuated by the absorption of gamma rays in the diffuse extragalactic background light through the QED interaction \(\gamma_{HE}\gamma_{EBL} \rightarrow f^+ f^-\). The cross section for this electromagnetic reaction decreases as the inverse of the square of the final state fermion mass and hence, the most probable final state is a \(e^+ e^-\) pair.

Gamma rays of energy \(E\) can interact with low-energy photons of energy \(\epsilon\) from the diffuse EBL over cosmological distance scales. The pair production is expected above the threshold energy condition

\[
E \epsilon (1 - \cos \theta) > 2m_e^2 c^4
\]

where \(\theta\) is the gamma-gamma scattering angle and \(m\) the fermion mass.

Therefore, the relevant EBL for the Cherenkov Telescopes is the visible and infra-red background, for which there exists observational data with determinations and bounds of the background spectral energy density (SED) at \(z = 0\) for several energies. The determinations come from direct measurements of the EBL density using instruments on satellites whereas the bounds, happen mostly in the infrared part of the EBL and come from extrapolations using galaxy counting. Given the difficulty of observing “cold galaxies” due to the zodiacal light background, they provide just lower limits.

Actually the SED at \(z = 0\) is not the end of the story since the EBL evolves with the redshift and the High-energy \(\gamma\)-rays originated at cosmological distances will interact with the EBL at different redshifts. The main contribution to the EBL comes from low-energy photons produced by stars in ordinary galaxies. Therefore either the star formation rate and the star evolution will play an important role to the EBL as a function of redshift determination.

The flux attenuation is a function of the gamma energy \(E\) and the redshift \(z\) of the gamma ray source and can be parameterised by the Optical depth \(\tau(E, z)\), which is defined as the number of e-fold reductions of the observed flux as compared with the initial flux at \(z\). This means that the Optical depth introduces an attenuation factor \(\exp[-\tau(E, z)]\) modifying the gamma ray source energy spectrum.

\[
\tau(E, z) = \int_0^z dt \, e^{\frac{E}{T(z')}} \int_0^{\frac{x}{2}} \frac{dx}{\sqrt{x}} \int_0^{\frac{x}{2m_e^2 c^4}} \frac{d\epsilon}{\epsilon} \cdot n(\epsilon, z') \cdot \sigma(2\epsilon x \epsilon(1 + z')^2)
\]

where \(n(\epsilon, z')\) is the EBL spectral density at redshift \(z'\), \(\sigma\) the cross-section for \(\gamma_{HE}\gamma_{EBL} \rightarrow e^+ e^-\) and \(dt/\sqrt{dx}\) the lookback time.

For any given gamma ray energy, the Gamma Ray Horizon (GRH) is defined as the source redshift for which the Optical depth is \(\tau(E, z) = 1\). Therefore, the GRH gives, for each gamma ray energy, the redshift location \(z\) of a source for which the intrinsic gamma flux suffers an e-fold decrease when observed on Earth \(z = 0\) due to the gamma-gamma absorption.

In practice, the cut-off due to the Optical depth is completely folded with the spectral emission of the gamma source. But on the other hand, the suppression factor in the gamma
Figure 4: Measured differential energy spectra of two of the farthest Blazars detected by HESS compared with plausible scenarios explored by HESS to explain the observed intrinsic spectrum of index 1.5.

 flux due to the Optical depth depends only (assuming a specific cosmology and spectral EBL density) on the gamma energy and the redshift of the source. Therefore, a common gamma energy spectrum behaviour of a set of different gamma sources at the same redshift is most likely due to the Optical Depth.

To compute the Optical depth using equation 2 there are two quantities which have to be known: on the one hand, the density of the EBL and its redshift dependence, and on the other hand, the cosmological evolution of our universe casted in the lookback time expression.

The direct measurement of the EBL density in the wavelength range relevant for VHE gamma ray absorption (from 0.1\(\mu m\) to 10\(\mu m\)) is very difficult because of our light-polluted environment, in particular by zodiacal light - sunlight reflected from dust clouds in our solar system. For this reason, the absorption measured by studying the distortion in the energy spectrum of distant sources, has already been widely used to try to bound the EBL density.

HESS\(^9\) and MAGIC\(^{10}\) have observed VHE gamma rays from few relatively distant active galaxies. In the case of HESS two objects, identified as the Blazars H2356-309 and 1ES1101-232 at redshifts of \(z = 0.165\) and 0.186 respectively, have been detected. The multiwavelength observations of Blazars as well as theoretical shock acceleration models in jets have serious difficulties to predict intrinsic gamma ray spectral energy slopes harder than \(\Gamma = 1.5\)\(^{9}\) while the observed slope for these two sources and for the 1ES1218+304 Blazar at redshift \(z = 0.182\) discovered by MAGIC are unexpectedly very hard, of about \(\Gamma \approx 3\) as can be seen in Figure 4. The observation of such hard spectra hints to a universe more transparent to VHE gamma rays than what was expected based on the direct measurements and the model predictions of the EBL density.

Actually, using these spectra and the energy dependence of the Optical depth through electron-positron pair production which can be obtained from equation 2, the HESS collaboration has been able to set a firm upper limit on the absorption of gamma ray and hence on the amount of extragalactic background light\(^{11}\).

This limit is sensibly less than - and hence in conflict with - the values derived by direct measurements of the extragalactic background light as can be seen in Figure 5. Furthermore, being only about a factor of \(\sim 1.5\) above the lower limit given by direct observation of galaxies by the Hubble Space Telescope, the HESS observations seriously limit the possible contribution

\(^{9}\)Nevertheless, it should be pointed out that this assumption could be relaxed in case of significant absorption of gamma rays at the source, for instance with the optical radiation from the accretion disk or scattered along the jet, which could produce an spectral index harder than 1.5.
from sources other than galaxies. This is in good agreement with recent theoretical calculations and arguments against a strong extragalactic background from first-generation stars. This is bad news for the attempts at direct detection of the glow of these population III stars but the HESS results expand the horizon of the gamma-ray universe, allowing Cherenkov telescopes to detect many other remote active galaxies.

The upper bound from HESS seems to be confirmed already by observations of new AGNs being recently reported by HESS and MAGIC. Therefore, taking into account that the correction of any possible observational biases in the galaxy count contribution to the EBL would very likely increase the lower bound, narrowing even further the distance between that lower bound and the HESS upper bound, one may think that the EBL density in the relevant region for VHE gamma ray astronomy might be basically resolved as the sum of the contributions from the light of all the galaxies observed as point-like sources. Since there are many deep-exposure large astronomical surveys in operation and proposed for the coming years cartographing the galaxies in big volumes of the visible universe, it may be then possible to get a rather accurate determination of the EBL density as a function of redshift in the wavelength region relevant for VHE gamma ray astronomy.

In that case, the only missing information in equation 2 would be the lookback time, and then the measurement of the Optical depth using the distant Blazar spectrum absorption could be turned upside down and used to try to measure the Cosmological Parameters instead of the EBL density.

Summarizing, there are two implications of the HESS results, namely:

- on the one hand, the universe is more transparent to gamma rays than expected and therefore the redshift reach of Cherenkov Telescopes should be substantially larger than anticipated allowing to observe much more distant extragalactic sources,

- on the other hand, the EBL density in the wavelength region relevant for the VHE gamma ray absorption might be actually resolved and hence the EBL density could be directly measured by surveys performing deep and detailed galaxy count

If these implications are confirmed, the study of the absorption in the energy spectrum of extragalactic VHE gamma rays at different redshifts may provide a competitive complementary technique for the determination of the parameters with govern the expansion of our universe and specifically, may help in constraining Dark Energy.\textsuperscript{12,13}

4 Tests of the invariance of the speed of light

Any quantum theory of gravitation introduces quantum fluctuations at the Planck scale \((E_P \approx 10^{13}\text{GeV or correspondingly } L_P \approx 10^{35}\text{cm})\), which would induce a deformed dispersion relation for photons of the form\textsuperscript{14}:

\[
p^2c^2 = E^2[1 + f(E/E_{QG})]
\]

where \(E\) is the photon energy, \(E_{QG}\) an effective quantum gravity energy scale (which might be as large as the Planck scale) and \(f\) is a model-dependent function of the ratio \(E/E_{QG}\), \(p\) is the photon momentum and \(c\) is the velocity of light. At small energies \(E \ll E_{QG}\) a series expansion of the dispersion relation can be made:

\[
p^2c^2 = E^2[1 + \xi E/E_{QG} + O(E^2/E_{QG}^2)]
\]

where \(\xi = \pm 1\) is a sign ambiguity which is fixed in the given theory. Equation 4 leads then to energy-dependent velocities of the photon:
\[ v = \frac{\partial E}{\partial p} \approx c \left( 1 - \frac{E}{E_{QG}} \right) \] (5)

Gamma rays travelling cosmological distances should therefore encounter a "vacuum" energy dispersion \( \delta v \sim E/E_{QG} \), violating Lorentz invariance. A gamma ray signal of observed energy \( E_0 \) should acquire a time delay with respect to the Lorentz-invariant case, after having travelled a distance \( L \) (redshift \( z \))\(^{15}\):

\[ \Delta t \approx \frac{E}{E_{QG}} \int_0^Z (1 + z) \frac{dl}{dz} dz \ll 1 \frac{E}{E_{QG} c} L \] (6)

Gamma rays of different energies being emitted simultaneously should thus reach an observer at different times. In order to use equation 6 to test \( E_{QG} \), a rapidly varying signal is required with typical time intervals \( \delta t \) smaller than the time delay \( \Delta t \) due to the quantum gravity effect and observed simultaneously at two different energies at least.

Gamma ray telescopes are specially well suited to measure this effect since they study photons at cosmological distances such as Blazars and Gamma Ray Bursts, and these sources provide natural time stamps since they are either flaring or transient. The light curves of these fast flares can be recorded and studied in detail thanks to the the huge effective areas of these telescopes.

Nevertheless, since possible energy-dependent time delays observed in a specific source could have an astrophysical origin and be produced either in the emission process or during the propagation of the photons thorough space for that specific region of the sky\(^{16}\), a sinequanone condition to make a claim of observation of a Quantum Gravity effect should be the observation of delays in a sample of sources distributed across different regions in the sky and located within a broad range of distances, which should nevertheless adjust the simple mathematical relation casted in equation 6.

In 1999, the Whipple collaboration published\(^{17}\) a first bound on \( E_{QG} \), obtained with that technique using a flare of Mrk 421 (\( z = 0.031 \)) which was very fast (\( \delta t \approx 280 \text{s} \) as can be seen in the lightcurve of figure 6) and was observed up to a gamma ray energy of 2 TeV. The analysis of that flare allowed the WHIPPLE collaboration to place a constraint of \( E_{QG} / \xi > 410^{16} \) GeV at 95% confidence level.

The MAGIC collaboration has recently reported\(^{10}\) recorded AGN flares from Mrk 501 even faster and with a much larger amount of gamma rays recorded than the one observed by WHIPPLE, allowing a broader and more detailed energy spectrum and which, therefore, may lead to much better bounds than the aforementioned one.
In addition, if GRB are detected with Cherenkov Telescopes, using the same method for GRBs, much higher sensitivities should be reached since the distances L are usually much larger and typical time intervals δt much shorter. For instance, assuming a GRB at a redshift of z = 1, observed simultaneously at 100 GeV and 1 MeV, with a time binning of 1 s, a hypothetical limit of $E_{\text{GRB}}/\xi > 10^{19}$ GeV could be reached. Therefore, IACTs might provide the opportunity of testing directly the quantum nature of Gravity up to effective scales of the order of the Planck mass.

References

XLIInd RENCONTRES DE MORIOND

Electroweak Interactions and Unified Theories

V - Baryo and Leptogenesis
RECENT ISSUES IN LEPTOGENESIS

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Baryogenesis via leptogenesis provides an appealing mechanism to explain the observed baryon asymmetry of the Universe. Recent refinements in the understanding of the dynamics of leptogenesis include detailed studies of the effects of lepton flavors and of the role possibly played by the lepton asymmetries generated in the decays of the heavier singlet neutrinos $N_{2,3}$. A review of these recent developments in the theory of leptogenesis is presented.

1 Introduction

The possibility that the Baryon Asymmetry of the Universe (BAU) could originate from a lepton number asymmetry generated in the $CP$ violating decays of the heavy seesaw Majorana neutrinos was put forth about twenty years ago by Fukugita and Yanagida. Their proposal came shortly after Kuzmin, Rubakov and Shaposhnikov pointed out that above the electroweak phase transition $B + L$ is violated by fast electroweak anomalous interactions. This implies that any lepton asymmetry generated in the unbroken phase would be unavoidably converted in part into a baryon asymmetry. However, the discovery that at $T \gtrsim 100$ GeV electroweak interactions do not conserve baryon number, also suggested the exciting possibility that baryogenesis could be a purely standard model (SM) phenomenon, and opened the way to electroweak baryogenesis. Even if rather soon it became clear that within the SM electroweak baryogenesis fails to reproduce the correct BAU by many orders of magnitude, within the minimal supersymmetric standard model (MSSM) the chances of success were much better, and this triggered an intense research activity in that direction. Indeed, in the early 90's electroweak baryogenesis attracted more interest than leptogenesis, but still a few remarkable papers appeared that put the first basis for quantitative studies of leptogenesis. Here I will just mention two important contributions that established the structure of the two main ingredients of leptogenesis: the rates for several washout processes relevant for the leptogenesis Boltzmann equations, that were presented by

Around year 2000 a flourishing of detailed studies of leptogenesis begins, with a corresponding burst in the number of papers dealing with this subject. This raise of interest in leptogenesis can be traced back to two main reasons: firstly, the experimental confirmation (from oscillation experiments) that neutrinos have nonvanishing masses strengthened the case for the seesaw mechanism, that in turn implies the existence, at some large energy scale, of lepton number violating ($L$) interactions. Secondly, the fact that the various analysis of supersymmetric electroweak baryogenesis cornered this possibility in a quite restricted region of parameter space, leaving for example for the Higgs mass just a 5 GeV window (115 - 120 GeV).

The number of important papers and the list of people that contributed to the development of leptogenesis studies and to understand the various implications for the low energy neutrino parameters is too large to be recalled here. However, let me mention the remarkable paper of Giudice et al. that appeared at the end of 2003: in this paper a whole set of thermal corrections for the relevant leptogenesis processes were carefully computed, a couple of mistakes common to previous studies were pointed out and corrected, and a detailed numerical analysis was presented both for the SM and the MSSM cases. Eventually, it was claimed that the residual numerical uncertainties would probably not exceed the 10%-20% level. A couple of years later, Nir, Roulet, Raaker and myself carried out a detailed study of additional effects that were not accounted for in the analysis of ref.. This included electroweak and QCD sphaleron effects, the effects of the asymmetry in the Higgs number density, as well as the constraints on the particles asymmetry-densities implied by the spectator reactions that are in thermal equilibrium in different temperature ranges relevant for leptogenesis. Indeed, we found that the largest of these new effects would barely reach the level of a few tens of percent.

However, two important ingredients had been overlooked in practically all previous studies, and had still to be accounted for. These were the role of the light lepton flavors, and the role of the heavier seesaw Majorana neutrinos. One remarkable exception was the 1999 paper by Barbieri et al. that, besides addressing as the main topic the issue of flavor effects in leptogenesis, also pointed out that the lepton number asymmetries generated in the decays of the heavier seesaw neutrinos can contribute to the final value of the BAU. However, these important results did not have much impact on subsequent analyses. The reason might be that these were thought to be just order one effects on the final value of the lepton asymmetry, with no other major consequences for leptogenesis. As I will discuss in the following, the size of the effects could easily reach the one order of magnitude level and, most importantly, they can spoil the leptogenesis constraints on the neutrino low energy parameters, and in particular the limit on the absolute scale of neutrino masses. This is important, since it was thought that this limit was a firm prediction of leptogenesis with hierarchical seesaw neutrinos, and that the discovery of a neutrino mass $m_\nu \gtrsim 0.2$ eV would have strongly disfavored leptogenesis, or hinted to different scenarios (as e.g. resonant leptogenesis).

2 The standard scenario

Let us start by writing the first few terms of the leptogenesis Lagrangian, neglecting for the moment the heavier neutrinos $N_{2,3}$ (except for their virtual effects in the $CP$ violating asymmetries):

$$\mathcal{L} = \frac{1}{2} \left[ \bar{N}_1 (i \not\partial N_1 - M_1 N_1 N_1) \right] - (\lambda_1 \bar{N}_1 \ell_i H + \text{h.c.}). \quad (1)$$

Lepton flavor effects were also considered by Endoh, Morozumi and Xiong in their 2003 paper, in the context of the minimal seesaw model with just two right handed neutrinos.
Here $N_1$ is the lightest right-handed Majorana neutrino with mass $M_1$, $H$ is the Higgs field, and $\ell_i$ is the lepton doublet to which $N_1$ couples, that when expressed on a complete orthogonal basis $\{\ell_i\}$ reads

$$|\ell_i\rangle = (\lambda \lambda^\dagger)_{11}^{-1/2} \sum_i \lambda_{iA} |\ell_i\rangle.$$

(2)

In practice it is always convenient to use the basis that diagonalizes the charged lepton Yukawa couplings (the flavor basis) that also has well defined CP conjugation properties $CP(\{\ell_i\}) = \{\bar{\ell}_i\}$ with $i = e, \mu, \tau$. Note that in the first and third term in (1) a lepton number can be assigned to $N_1$, that is however violated by two units by the mass term. Then eq. (1) implies processes that violate $L$, like inverse-decays followed by $N_1$ decays $\ell_1 \leftrightarrow N_1 \leftrightarrow \ell_1$, off-shell $\Delta L = 2$ scatterings $\ell_1 H \leftrightarrow \ell_1 \bar{H}$, $\Delta L = 1$ scatterings involving the top-quark like $N_1 \ell_1 \leftrightarrow Q_3 l$ or involving the gauge bosons like $N_1 \ell_1 \rightarrow AH$ (with $A = W, B$). The temperature range in which $\ell$ processes can be important for leptogenesis is around $T \sim M_1$. This is because if the $\lambda_1$ couplings were large enough that these processes were already relevant at $T \gg M_1$ (when the Universe expansion is fast) than they would come into complete thermal equilibrium at lower temperatures (when the expansion slows down) thus forbidding the survival of any macroscopic $L$ asymmetry. On the other hand at $T \ll M_1$ decays, inverse decays and $\Delta L = 1$ scatterings are Boltzmann suppressed, $\Delta L = 2$ scatterings are power suppressed, and therefore $L$ violating processes become quite inefficient as the temperature drops well below $M_1$.

The possibility of generating an asymmetric between the number of leptons $n_{\ell_1}$ and antileptons $n_{\bar{\ell}_1}$ is due to a non-vanishing CP asymmetry in $N_1$ decays:

$$\epsilon_1 \equiv \frac{\Gamma(N_1 \rightarrow \ell_1 H) - \Gamma(N_1 \rightarrow \ell_1 \bar{H})}{\Gamma(N_1 \rightarrow \ell_1 H) + \Gamma(N_1 \rightarrow \ell_1 \bar{H})} \neq 0.$$

(3)

In order that a macroscopic $L$ asymmetry can build up, the condition that $L$ reactions are (at least slightly) out of equilibrium at $T \sim M_1$ must also be satisfied. This condition can be expressed in terms of two dimensionless parameters, defined in terms of the Higgs vev $v \equiv \langle H \rangle$ and of the Plank mass $M_P$ as:

$$\tilde{m}_1 = \frac{(\lambda \lambda^\dagger)_{11} v^2}{M_1}, \quad m_* \approx 10^3 \frac{\langle H \rangle^2}{M_P} \approx 10^{-3} \text{ eV}.$$

(4)

The first parameter ($\tilde{m}_1$) is related to the rates of $N_1$ processes (like decays and inverse decays) while the second one ($m_*$) is related to the expansion rate of the Universe at $T \sim M_1$. When $\tilde{m}_1 \lesssim m_*$, $L$ processes are slower than the Universe expansion rate and leptogenesis can occur. As $\tilde{m}_1$ increases to values larger than $m_*$, $L$ reactions approach thermal equilibrium thus rendering leptogenesis inefficient because of the back-reactions that tend to erase any macroscopic asymmetry. However, even for $\tilde{m}_1$ as large as $\sim 100 m_*$ a lepton asymmetry sufficient to explain the BAU can be generated. It is customary to refer to the condition $\tilde{m}_1 > m_*$ as to the strong washout regime since washout reactions are rather fast. This regime is considered more likely than the weak washout regime $\tilde{m}_1 < m_*$ in view of the experimental values of the light neutrino mass-squared differences (that are both $> m_\nu^2$) and of the theoretical lower bound $\tilde{m}_1 \geq m_{\nu_3}$, where $m_{\nu_3}$ is the mass of the lightest neutrino. The strong washout regime is also theoretically more appealing since the final value of the lepton asymmetry is independent of the particular value of the $N_1$ initial abundance, and also of a possible asymmetry $Y_{\ell_1} = (n_{\ell_1} - n_{\bar{\ell}_1})/s \neq 0$ (where $s$ is the entropy density) preexisting the $N_1$ decay era. This last fact has been often used to argue that for $\tilde{m}_1 > m_*$ only the dynamics of the lightest Majorana neutrino $N_1$ is important, since asymmetries generated in the decays of the heavier $N_{2,3}$ would be efficiently erased by the strong $N_1$-related washouts. As we will see below, the effects of $N_1$ interactions on the $Y_{\ell_{2,3}}$ asymmetries are subtle, and the previous argument is incorrect. The result of numerical
integration of the Boltzmann equations for $Y_{\ell_1}$ can be conveniently expressed in terms of an efficiency factor $\eta_1$, that ranges between 0 and 1:

$$Y_{\ell_1} = 3.9 \times 10^{-3} \eta_1 \varepsilon_1,$$

$$\eta_1 \approx \frac{m_\star}{m_1}. \tag{5}$$

The second relation gives a rough approximation for $\eta_1$ in the strong washout regime, that will become useful in analyzing the impact of flavor effects. Clearly, too strong washouts ($m_1 \gg m_\star$) can put in jeopardy the success of leptogenesis by suppressing too much the efficiency. However, it should also be stressed that washouts constitute a fundamental ingredient to generate a lepton asymmetry. This is particularly true in thermal leptogenesis, with zero initial $N_1$ abundance, and is illustrated in fig. 1 where the evolution of the lepton asymmetry for a representative model is plotted against decreasing values of the temperature. The different curves correspond to different level of (artificial) reduction in the strength of the washout rates (but not in the $N_1$ production rates) from the model value (solid red line), to 10% (dashed blue line), 1% (dot-dashed pink line) and 0.1% (dotted green line). The solid black line corresponds to switching off all back-reactions. (Of course the last four curves correspond to unphysical conditions.) It is apparent that while a partial reduction in the washout rates is beneficial to leptogenesis, an excessive reduction suppresses the final asymmetry and eventually, when washouts are switched off completely, no asymmetry survives. This behavior can be understood as follows: all leptogenesis processes can be seen as scatterings between standard model particle states $X, Y$ involving intermediate on-shell and off-shell unstable $N_1$'s: $X \leftrightarrow N_1^{(s)} \leftrightarrow Y$. Since the CP asymmetry of any $X \leftrightarrow Y$ process is at most of $O(\lambda_1^5)^{15}$, if the lepton asymmetries generated in the different processes were exactly conserved, the overall amount that could survive would not exceed this order. Moreover, since $O(\lambda_1^5)$ asymmetries are systematically neglected in the Boltzmann equations, the numerical result would be exactly zero. However, the on-shell and off-shell components of each process have much larger CP asymmetries of $O(\lambda_1^4)$, and the cancellation to $O(\lambda_1^5)$ occurs because they are
of opposite sign and (at leading order in the couplings) of the same magnitude. Moreover, since
the long range and short range components of each process have different time scales, at each
instant during leptogenesis a lepton asymmetry up to $\mathcal{O}(\lambda_1^4)$ can be present. Washout
processes by definition do not conserve the lepton asymmetries, and most importantly they act unevenly
over the different processes as well as over their short and long range components, erasing more
efficiently the asymmetries generated in $N_1$ production processes and off-shell scatterings that on
average occur at earlier times, and washing out less efficiently the asymmetries of processes that
destroy $N_1$'s (on-shell scatterings and decays). It is thanks to the washouts that an unbalanced
lepton asymmetry up to $\mathcal{O}(\lambda_1^4)$ can eventually survive. In the next section we will see that when
flavor effects are important, washouts can play an even more dramatic role in leptogenesis.

The possibility of deriving an upper limit for the the light neutrino masses\textsuperscript{13} follows from
the existence of a theoretical bound on the maximum value of the $CP$ asymmetry $\epsilon_1$ (that holds
when $N_{1,2,3}$ are sufficiently hierarchical, and $m_{\nu_{1,2,3}}$ quasi degenerate) and relates $M_1$, $m_{\nu_3}$ and
the washout parameter $\tilde{m}_1$:

$$|\epsilon_1| \leq \left[ \frac{3}{16\pi} \frac{M_1}{v^2} \left( m_{\nu_3} - m_{\nu_1} \right) \right] \sqrt{1 - \frac{m_{\nu_3}^2}{\tilde{m}_1^2}}. \quad (6)$$

The term in square brackets is the so called Davidson-Ibarra limit\textsuperscript{15} while the correction in the
square root was first given in ref.\textsuperscript{17} When $m_{\nu_3} \gtrsim 0.1$ eV, the light neutrinos are quasi-degenerate
and $m_{\nu_3} - m_{\nu_1} \sim \Delta m_{\text{atm}}^2 / 2 m_{\nu_3} \rightarrow 0$ so that, to keep $\epsilon_1$ finite, $M_1$ is pushed to large values
$\gtrsim 10^{15}$ GeV. Since at the same time $\tilde{m}_1$ must remain larger than $m_{\nu_3}$, the washouts also increase,
until the surviving asymmetry is too small to explain the BAU.\textsuperscript{5} The limit $m_{\nu_3} \lesssim 0.15$ eV results.

3 Lepton flavor effects

In the Lagrangian (1) the terms involving the charged lepton Yukawa couplings have not been
included. Since all these couplings are rather small, if leptogenesis occurs at temperatures $T \gtrsim 10^{12}$ GeV, when the Universe is still very young, not many of the related (slow) processes
could have occurred during its short lifetime, and leptogenesis has essentially no knowledge of
lepton flavors. At $T \lesssim 10^{12}$ GeV the reactions mediated by the tau Yukawa coupling $h_\tau$
become important, and at $T \lesssim 10^9$ GeV also $h_\mu$-reactions have to be accounted for. Including
the Yukawa terms for the leptons yields the Lagrangian:

$$\mathcal{L} = \frac{1}{2} \left[ N_1 (i \partial \bar{\psi}) N_1 - M_1 N_1 N_1 \right] - (\lambda_{14} N_1 \ell_i H + h_{i\ell} \bar{\ell}_i \ell H^\dagger + \text{h.c.}), \quad (7)$$

where (in the flavor basis) the matrix $h$ of the Yukawa couplings is diagonal. The flavor content
of the (anti)lepton doublets $\ell_i$ ($\bar{\ell}_i$) to which $N_1$ decays is now important, since these states do
not remain coherent, but are effectively resolved into their flavor components by the fast Yukawa
interactions $h_i$.\textsuperscript{11,18,19} Note that because of $CP$ violating loop effects, in general $CP(\bar{\ell}_i) \neq \ell_i$,
that is the antileptons produced in $N_1$ decays are not the $CP$ conjugate of the leptons, implying
that the flavor projections $K_i \equiv |\langle \ell_i | \ell_i \rangle|^2$ and $\bar{K}_i \equiv |\langle \bar{\ell}_i | \bar{\ell}_i \rangle|^2$ differ: $\Delta K_i = K_i - \bar{K}_i \neq 0$. The flavor $CP$ asymmetries can be defined as:\textsuperscript{13}

$$\epsilon_i' = \frac{\Gamma(N_1 \rightarrow \ell_i H) - \Gamma(N_1 \rightarrow \bar{\ell}_i \bar{H})}{\Gamma_{N_1}} = K_i \epsilon_1 + \Delta K_i / 2. \quad (8)$$

The factor $\Delta K_i$ in the second equality accounts for the flavor mismatch between leptons and
antileptons. The factor $K_i$ in front of $\epsilon_1$ accounts for the reduction in the strength of the $N_1-\ell_i$
coupling with respect to $N_1$-$\ell_1$, and thus reduces also the strength of the washouts for the $i$-flavor, yielding an efficiency factor $\eta_i^d = \min(\eta_i/K_i, 1)$. Assuming for illustration $\eta_i/K_i < 1$ the resulting asymmetry is

$$Y_L \approx \sum_i \epsilon_i \eta_i^d \approx n_f Y_{\ell_1} + \sum_i \frac{\Delta K_i}{2 K_i} m_\ell \frac{m_\ell}{m_1}.$$  \hspace{1cm} (9)

In the first term on the r.h.s. $n_f$ represents the number of flavors effectively resolved by the charged lepton Yukawa interactions ($n_f = 2$ or 3) while $Y_{\ell_1}$ is the asymmetry that would have been obtained by neglecting the decoherence of $\ell_1$. The second term, that is controlled by the ‘flavor mismatch’ factor $\Delta K_i$, can become particularly large in the cases when the flavor $i$ is almost decoupled from $N_1$ ($K_i \ll 1$). This situation is depicted in fig. 2 for the two-flavor case and for two different temperature regimes. The two flat curves give $|Y_{B-L}|$ as a function of the flavor projector $K_\ell$ assuming $\Delta K_\ell = 0$, and show rather clearly the enhancement of a factor $\approx 2$ with respect to the one flavor case (the points at $K_\ell = 0, 1$). The other two curves are peaked at values close to the boundaries, when $\ell_\tau$ or a combination orthogonal to $\ell_\tau$ are almost decoupled from $N_1$, and show how the $\ell_1$-$\ell_1$ flavor mismatch can produce much larger enhancements.

It was first noted in ref.\textsuperscript{15} that flavored-leptogenesis can be viable even when the branching ratios for decays into leptons and antileptons are equal, that is in the limit when $L$ is conserved in decays and the total asymmetry $\epsilon_1 = 0$ vanishes. This is a surprising possibility, that can occur when the $CP$ asymmetries for the single flavors are non-vanishing, thanks to the fact that lepton number is in any case violated by the washout interactions.

In conclusion, the relevance of flavor effects is at least twofold:

1. The BAU resulting from leptogenesis can be several times larger than what would be obtained neglecting flavor effects.

2. If leptogenesis occurs at temperatures when flavor effects are important, the limit on the light neutrino masses does not hold.\textsuperscript{18,20} This is because there is no analogous of the Davidson-Ibarra bound in eq. (6) for the flavor asymmetries $\epsilon_i$. 
The effects of the heavier Majorana neutrinos

What about the possible effects of the heavier Majorana neutrinos $N_{2,3}$ that we have so far neglected? A few recent studies analyzed the so called \textquoteleft\textquoteleft N$_1$-decoupling\textquoteright\textquoteright scenario, in which the Yukawa couplings of $N_1$ are simply too weak to washout the asymmetry generated in $N_2$ decays (and $\epsilon_1$ is too small to explain the BAU).\cite{21} This is a consistent scenario in which $N_2$ leptogenesis could successfully generate enough lepton asymmetry. However, in the opposite situation when the Yukawa couplings of $N_1$ are very large, it was generally assumed that the asymmetries related to $N_{2,3}$ are irrelevant, since they would be washed out during $N_1$ leptogenesis. In contrast to this, in ref.\cite{11} (and more recently also in ref.\cite{22}) it was stated that part of the asymmetry from $N_{2,3}$ decays does in general survive, and must be taken into account when computing the BAU.

In ref.\cite{23} Engelhard, Grossman, Nir and myself carried out a detailed study of the fate of a lepton asymmetry $Y_P$ preexisting $N_1$ leptogenesis, and we reached conclusions that agree with these statements. I will briefly describe the reasons for this and the importance of the results. Including also $N_{2,3}$ the leptogenesis Lagrangian reads:

$$\mathcal{L} = \frac{1}{2} [\bar{N}_\alpha (i \gamma^\mu) N_\alpha - N_\alpha M_\alpha N_\alpha] - (\lambda_{\alpha i} \bar{N}_\alpha \ell_i H + h.c.),$$  \hspace{1cm} (10)

where the heavy neutrinos are written in the mass basis with $\alpha = 1, 2, 3$. It is convenient to define the three (in general non-orthogonal) combinations of lepton doublets $\ell_\alpha$ to which the corresponding $N_\alpha$ decay:

$$|\ell_\alpha\rangle = (\lambda \lambda^\dagger)^{-1/2}_{\alpha i} \sum_i \lambda_{\alpha i} |\ell_i\rangle.$$  \hspace{1cm} (11)

Let us discuss for definiteness the case when $N_2$-related washouts are not too strong ($\bar{m}_2 \gg m_\alpha$)\hspace{1cm}, so that a sizeable asymmetry proportional to $\epsilon_2$ is generated, while $N_1$-related washouts are so strong that by itself $N_1$ leptogenesis would not be successful ($\bar{m}_1 \gg m_\alpha$). To simplify the arguments, let us also impose two additional conditions: thermal leptogenesis, that is a vanishing initial $N_1$ abundance $n_{N_1} (T \gg M_1) \approx 0$, and a strong hierarchy $M_2 / M_1 \gg 1$. From this it follows that there are no $N_1$ related washout effects during $N_2$ leptogenesis and, because $n_{N_2} (T \approx M_1)$ is Boltzmann suppressed, there are no $N_2$ related washouts during $N_1$ leptogenesis. Thus $N_2$ and $N_1$ dynamics are decoupled. Now, the second condition in (??) implies that already at $T \gtrsim M_1$ the interactions mediated by the $N_1$ Yukawa couplings are sufficiently fast to quickly destroy the coherence of $\ell_2$ produced in $N_2$ decays. Then a statistical mixture of $\ell_1$ and of states orthogonal to $\ell_1$ builds up, and it can be described by a suitable \emph{diagonal} density matrix. For simplicity, let us assume that both $N_2$ and $N_1$ decay at $T \gtrsim 10^{12}$ GeV when flavor effects are irrelevant. In this regime a convenient choice for the orthogonal lepton basis is $(\ell_{i_1}, \ell_{i_1 \perp})$ where $\ell_{i_1 \perp}$ denotes generically the flavor components orthogonal to $\ell_{i_1}$. Then any asymmetry $Y_P$ preexisting the $N_1$ leptogenesis phase (as for example $Y_{\ell_{i_1}}$) decomposes as:

$$Y_P = Y_{\ell_{i_1 \perp}} + Y_{\ell_{i_1}}.$$  \hspace{1cm} (12)

The crucial point here is that in general we can expect $Y_{\ell_{i_1 \perp}}$ to be of the same order than $Y_P$, and since $\ell_{i_1 \perp}$ is orthogonal to $\ell_{i_1}$, this component of the asymmetry remains protected against $N_1$ washouts. Therefore, a finite part of any preexisting asymmetry (and in particular of $Y_{\ell_{i_1}}$ generated in $N_2$ decays) survives through $N_1$ leptogenesis. A more detailed study\cite{23} reveals also some additional features. For example, in spite of the strong $N_1$-related washouts $Y_{\ell_{i_1}}$ is not driven to zero, rather, only the sum of $Y_{\ell_{i_1}}$ and of the Higgs asymmetry $Y_H$ vanishes, but not the two separately. (This can be traced back to the presence of a conserved charge related to $Y_{\ell_{i_1}}$.)

For $10^9 \lesssim M_1 \lesssim 10^{12}$ GeV the lepton flavor structures are only partially resolved during $N_1$ leptogenesis, and a similar result is obtained. However, when $M_1 \lesssim 10^9$ GeV and the full flavor basis $(\ell_e, \ell_\mu, \ell_\tau)$ is resolved, there are no directions in flavor space where an asymmetry can
remain protected, and then $Y_P$ can be completely erased independently of its flavor composition. In conclusion, the common assumption that when $N_1$ leptogenesis occurs in the strong washout regime the final BAU is independent of initial conditions, does not hold in general, and is justified only in the following cases:\textsuperscript{23} \textit{i)} Vanishing decay asymmetries and/or efficiency factors for $N_{2,3}$ ($\varepsilon_{2\eta_2} \approx 0$ and $\varepsilon_{3\eta_3} \approx 0$); \textit{ii)} $N_1$-related washouts are still significant at $T \lesssim 10^9$ GeV; \textit{iii)} Reheating occurs at a temperature in between $M_2$ and $M_1$. In all other cases the initial conditions for $N_1$ leptogenesis, and in particular those related to the possible presence of an initial asymmetry from $N_{2,3}$ decays, cannot be ignored when calculating the BAU, and any constraint inferred from analyses based only on $N_1$ leptogenesis are not reliable.

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Minimal Lepton Flavour Violation and leptogenesis

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We investigate the viability of leptogenesis in models with three heavy right-handed neutrinos, where the charged-lepton and the neutrino Yukawa couplings are the only irreducible sources of lepton-flavour symmetry breaking (Minimal Lepton Flavour Violation hypothesis). We analyze the impact of the CP violating phases responsible for leptogenesis on low-energy flavour changing observables.

1 Introduction

In the Standard Model (SM) the structure of the fermion mass matrices controls the physics of flavour. In the quark sector for instance, up and down quarks have mass matrices which are approximately aligned and the eigenvalues are several orders of magnitude smaller than the electroweak scale. This results in effective GIM and Cabibbo suppressions of flavour changing rates and guarantees a good level of predictivity to the theory. It would be useful to extend the SM connection between flavour and masses in models of new physics which predict new flavoured particles at the TeV scale, that could easily conflict with bounds from flavour experiments. This is the idea behind the Minimal Flavour Violation (MFV) principle\textsuperscript{1,2,3}. It is a general and flexible hypothesis that has been implemented in strongly-interacting theories\textsuperscript{4}, low-energy supersymmetry\textsuperscript{2,3}, multi Higgs\textsuperscript{3,4} and GUT\textsuperscript{5} models.

The SM Lagrangian is modified by the presence of high-energy new physics by the appearance of non-renormalizable operators, constructed from SM and Yukawa fields and suppressed by inverse powers of the mass of the new flavoured particles, $A_{\text{LFV}}$. The MFV recipe for model building requires that the new physics operators obey those flavour symmetry properties of the SM which are responsible for the suppression of flavour changing effects. In the quark sector there is essentially no ambiguity because the largest group of unitarity field transformations that commutes with the gauge group, $\mathcal{G}_q = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_Y$, is broken only by the Yukawa couplings. The invariance of the SM Lagrangian under $\mathcal{G}_q$ can be formally
recovered elevating the Yukawa matrices to spurion fields with the appropriate transformation properties under $G_q$. The hypothesis of MFV states that this must be sufficient to make invariant under $G_q$ also the new physics operators. The quark sector has been extensively tested in flavour experiments and it seems plausible that some kind of MFV principle is really at work in the underlying theory.

Apart from the analogy with quarks, the definition of a Minimal Lepton Flavour Violation (MLFV) principle is demanded by a severe flavour problem in the lepton sector as well. For instance, the charged lepton radiative decays $l_i \rightarrow l_j \gamma$ are mediated by the effective operator

$$
\frac{\Delta_{i\beta}^{RL}}{\Lambda_{LFV}^2} g H^\dagger \epsilon_R \sigma^\beta \epsilon^\dagger L_i F_{\sigma\rho} = \frac{\Delta_{i\beta}^{RL}}{\Lambda_{LFV}^2} O_{RL}
$$

(1)

where the matrix $F_{\sigma\rho}$ is the field strength tensor of the gauge group, $g$ the corresponding gauge coupling and $\Delta_{i\beta}^{RL}$ is a flavour changing matrix which mediates the transition. If unnaturally small dimensionless parameters do not enter this expression, the experimental limit for $\mu \rightarrow e \gamma$ pushes the scale $\Lambda_{LFV}$ to values $>10^5$ TeV, well above the electroweak scale.

The implementation of a MFV principle in the lepton sector is not as straightforward as in the quark sector. If neutrinos are Majorana particles, neutrino masses themselves are generated by a non-renormalizable operator suppressed by the scale of lepton number violation (LNV)

$$
m_{\nu} \propto g_{\nu} \left( H L L \right)_i \left( H L L \right)_j
\frac{1}{\Lambda_{LNV}}
$$

(2)

so that (i) we face a two scale problem, probably with the hierarchy $\Lambda_{LNV} \gg \Lambda_{LFV}$; (ii) and we have to make hypotheses about the origin of the effective operator in eq.(2). In this work, we assume that neutrino masses are generated by the standard see-saw mechanism with three heavy right-handed neutrinos and we explore the connection with the generation of the baryon asymmetry of the universe. However different implementations of the MFV idea in the lepton sector are possible and do not necessarily involve an extension of the SM field content.

2 MLFV with three right-handed neutrinos

We consider three right-handed Majorana neutrinos $\nu^R_i$ in addition to the three SM left-handed doublets $L^c_i$ and right-handed charged lepton singlets $\epsilon^c_R$. The maximal flavour group is

$$
G_L = SU(3)_{L_L} \times SU(3)_{e_R} \times O(3)_{\nu_R},
$$

(3)

in close analogy with quark sector: the $\nu^R_i$'s are the counterpart of right-handed up quarks and the irreducible sources of flavour breaking are the charged lepton and neutrino Yukawa couplings. We choose not to take the Majorana mass term for right-handed neutrinos as an independent spurion, but we assume that it is proportional to the identity matrix in flavour space,

$$
M_R = \Lambda_{LNV} \mathbf{I}
$$

(4)

thus breaking the $SU(3)_{\nu_R}$ symmetry to $O(3)_{\nu_R}$. This assumption increases the predictivity of the model and can be excluded experimentally.

The symmetry breaking lagrangian is

$$
L_{Sym.Br.} = -\lambda^{ij}_L \bar{e}^c_R (H^\dagger L_i^c) + i \lambda^{ij}_Y \bar{\nu}^c_R (H^T \tau_j L_i^c) + \text{h.c.}
$$

$$
\rightarrow -v \lambda^{ij}_L \bar{e}^c_R \epsilon^c_R - v \lambda^{ij}_Y \bar{\nu}^c_R \nu^c_R + \text{h.c.}
$$

(5)

The irreducible sources of lepton symmetry breaking are the Yukawa matrices $\lambda_L$ and $\lambda_Y$, transforming as $(3,1,3)$ and $(3,3,1)$ under the symmetry group $G_L$. 
3 Connection between low energy observables and leptogenesis

Having identified the irreducible sources of flavour symmetry breaking and their transformation properties under the flavour group $G_f$, we are now ready to build the non-renormalizable operators suppressed by inverse powers of $\Lambda_{LFV}$ that contribute to flavour violating processes. According to the MLFV hypothesis, these operators must be invariant combinations of the SM fields and the spurions $\lambda_\nu$ and $\bar{\lambda}_\nu$. The complete list of the operators contributing at leading order to LFV processes is given in Refs.\textsuperscript{6,7}. In the case of charged lepton radiative decays there are two six-dimensional operators,

$$O_{RL}^{(1)} = g' H^\dagger e_R \sigma^{\mu\nu} \lambda_\nu \Delta_{RL}^{\nu L} B_{\mu\nu},$$
$$O_{RL}^{(2)} = g H^\dagger e_R \sigma^{\mu\nu} \lambda_\nu \Delta_{FCNC}^{\nu L} W_{\mu\nu}^a,$$  

which enter the effective Lagrangian for $l_i \rightarrow l_j \gamma$:

$$\mathcal{L}_{\text{eff}} = \frac{1}{\Lambda_{LFV}^2} \left( c_{RL}^{(1)} O_{RL}^{(1)} + c_{RL}^{(2)} O_{RL}^{(2)} \right),$$

where $g'$ ($g$) and $B_{\mu\nu}$ ($W_{\mu\nu}^a$) are the coupling constant and the field strength tensor of the $U(1)_Y$ ($SU(2)_L$) gauge group.

The matrix $\Delta_{RL}$ is a spurion combination transforming as $\bar{(3)}$ under the flavour group $G_f$. In the basis where $\lambda_\nu$ is diagonal and at leading order in $\lambda_\nu = m_\nu/v$, it is given by

$$\Delta_{RL} = \lambda_\nu (\lambda_\nu^T \lambda_\nu)$$

In the quark sector the analogue quantity $\lambda_q^T \lambda_q$ can be expressed in terms of quark masses and CKM angles, so that the only unknown in the predictions for flavour changing branching ratios is the overall normalization given by $\Lambda_{LFV}$, coming from the power suppression of new physics operators:

$$BR(\text{quark MFV}) \sim \frac{f(m_u, m_d, V_{CKM})}{\Lambda_{LFV}^2}$$

In the lepton sector instead we can obtain $\lambda_\nu^T \lambda_\nu$ extracting $\lambda_\nu$ from the see-saw relation

$$m_{\nu} \equiv v^2 \lambda_\nu^T M_R^{-1} \lambda_\nu = U_{PMNS}^\dagger m_{\text{diag}} U_{PMNS}^\dagger$$

written in the basis where the charged lepton Yukawa matrix $\lambda_\nu$ and the heavy Majorana neutrinos mass matrix $M_R$ are diagonal. We find

$$\lambda_\nu = \frac{M_R^{1/2}}{v} \lambda_\nu [H(\phi_1, \phi_2, \phi_3) m_{\text{diag}}^{1/2} U_{PMNS}^\dagger$$

where $O$ is a real orthogonal matrix which depends on three real parameters and $H$ a complex hermitian matrix which also depends on three real parameters $\phi_{1,2,3}$; $H \rightarrow I$ in the CP limit.\textsuperscript{11} In our definition of MFV, right-handed neutrinos are degenerate in mass so that $M_R$ in (11) is a number that we identify with $\Lambda_{LFV}$ and the matrix $O$ can be rotated away (the Lagrangian is $O(3)$ invariant). Using eq.(11), the basic unit that takes part into low energy flavour changing rates in the MFV framework can be written as

$$\lambda_\nu^T \lambda_\nu = \frac{M_R}{v^2} U_{PMNS}^\dagger m_{\text{diag}}^{1/2} H^{1/2} m_{\text{diag}}^{1/2} U_{PMNS}^\dagger.$$

Let us count the parameters in eq.(12): there are 9 parameters in principle measurable at low energy (MNS angles, Dirac and Majorana phases, neutrino masses) and 4 unknown parameters:
the normalization $M_R$ and $\phi_{1,2,3}$ in the matrix $H$, which disappear in the see-saw relation (10). With non-degenerate right-handed neutrinos we would have 5 unknown parameters more (the 2 mass splittings of $M_R$ and the 3 angles in $O$).

Assuming that all the baryon asymmetry $\eta_B$ of the universe has been generated through sphaleron effects by the lepton asymmetry produced in the out-of-equilibrium decays of right-handed neutrinos, we can use the observed value of $\eta_B = (6.3 \pm 0.3) \times 10^{-10}$ to get some information on the high energy parameters $M_R$ and $\phi_{1,2,3}$. In fact, in the one-flavour approximation leptogenesis depends on the combination

$$\lambda_\nu \lambda_\nu^\dagger = \frac{M_R}{v^2} H m_{\text{diag}} H.$$  

(13)

4 Analysis of leptogenesis

A necessary condition for generating a lepton asymmetry is the non-degeneracy of heavy neutrinos. We generally expect that the tree-level degeneracy of heavy neutrinos is lifted by radiative corrections. In MFV models the most general form of the allowed mass-splittings is

$$\frac{\Delta M_R}{M_R} = c_\nu \left[ \lambda_\nu \lambda_\nu^\dagger + (\lambda_\nu \lambda_\nu^\dagger)^T \right]$$
$$+ \ c^{(1)}_{\nu \nu} \left[ \lambda_\nu \lambda_\nu^\dagger \lambda_\nu \lambda_\nu^\dagger + (\lambda_\nu \lambda_\nu^\dagger \lambda_\nu \lambda_\nu^\dagger)^T \right]$$
$$+ \ c^{(2)}_{\nu \nu} \left[ (\lambda_\nu \lambda_\nu^\dagger)^T \lambda_\nu \lambda_\nu^\dagger \right]$$
$$+ \ c_{\text{det}} \left[ \lambda_\nu \lambda_\nu^\dagger \lambda_\nu \lambda_\nu^\dagger + (\lambda_\nu \lambda_\nu^\dagger \lambda_\nu \lambda_\nu^\dagger)^T \right] + \ldots$$

(14)

In a model-independent approach we cannot fix the coefficients in (14) but in a perturbative scenario we expect $c_\nu \sim g_{\text{eff}}^2 / 4\pi$ and $c^{(1)}_{\nu \nu}, c^{(2)}_{\nu \nu} \sim g_{\text{eff}}^2 / (4\pi)^2$. From eq.(14) we can derive some general properties of leptogenesis in MFV models: (i) the term proportional to $c_\nu$ does not generate an asymmetry by itself, but (ii) sets the order of magnitude of the mass splitting and naturally gives the condition for resonant leptogenesis: the mass splitting of right-handed neutrinos is comparable to the decay width,

$$\Delta M_{R,i} \sim \Gamma_i = M_R \frac{\left| \lambda_\nu \lambda_\nu^\dagger \right|^2 \Gamma_j}{8\pi}.$$  

(15)

(iii) The right amount of leptogenesis can be generated even with $\lambda_\nu = 0$, provided that all the three parameters $\phi_{1,2,3} \neq 0$. Since $\lambda_\nu \sim \sqrt{M_R}$, (iv) for low values of $M_R \lesssim 10^{12}$ GeV the asymmetry generated by the $c_{\text{det}}$ term dominates but is typically too small to match the observed value of $\eta_B$. In this regime we find the flat dependence on $M_R$ typical of resonant leptogenesis. At $M_R \gtrsim 10^{12}$ GeV the quadratic terms $c^{(1)}_{\nu \nu}$ dominate the generation of the asymmetry, which grows linearly with $M_R$. These specific features of resonant leptogenesis in MFV can be derived with a general analysis of CP-invariants and reproduced analytically\(^7\). Properties (i),(iv) explain the characteristic behaviour of $\eta_B$ as function of $M_R$, shown in fig. 1.

In deriving this result we used the analytic formulae for leptogenesis without flavour effects of Ref.\(^{12}\) and assumed a loop hierarchy between the coefficients of the mass splittings. Under these assumptions, we find that the right size for $\eta_B$ can be reached for $M_R \gtrsim 10^{12}$ GeV. The regime $M_R \gg 10^{12}$ GeV is particularly interesting since in this case the CP-violating parameters $\phi_0$ are very small and we recover the predictive scheme of Ref.\(^6\).

A MFV model is for instance the Minimal Supersymmetric Standard Model with degenerate right-handed neutrinos at the GUT scale. This scenario has been analyzed in Ref.\(^{10}\) where also flavour effects were included in the leptogenesis analysis.
5 Results for lepton radiative decays

In the MLFV framework, the effective Lagrangian in eq.(2) relevant for lepton radiative decays leads to the branching ratio \( l_i \to l_j \gamma \)

\[
\text{BR}_{l_i \to l_j \gamma} = \frac{\Gamma(l_i \to l_j \gamma)}{\Gamma(l_i \to l_j \nu \bar{\nu})} = 384 \pi^2 e^2 \frac{\alpha}{\Lambda_{\text{MFV}}^2} \left| \left( \lambda^+ \lambda^- \right)_{ij} \right|^2 \left| \epsilon^{(2)}_{\text{RL}} - \epsilon^{(1)}_{\text{RL}} \right|^2
\]

\[
\sim \left( \frac{\Lambda_{\text{LNV}}}{\Lambda_{\text{MFV}}} \right)^2 f \left( m_{\nu}^2, m_t, U_{PMNS}, \phi_1 \right)
\]

(16)

to be compared with eq.(9), which represents a generic branching ratio in MFV in the quark sector. An important difference is that in the lepton sector the \( \phi_i \) parameters, non-measurable in low energy experiments, enter the branching ratio. As shown in fig. 2, they typically produce an enhancement, but weaken the predictivity of the model. However for large values of \( M_R \) their effect is moderate and MFV predicts \( \text{BR}(\mu \to e\gamma)/\text{BR}(\tau \to \mu\gamma) \leq 1 \). The second important difference between eqs.(9) and (16) is the absolute normalization of the branching ratios, that now depends on both the scales of lepton number and lepton flavour violation. For this reason, we cannot directly interpret the experimental bounds on \( l_i \to l_j \gamma \) as lower limits on the scale \( \Lambda_{\text{FV}} \) without independent information on \( \Lambda_{\text{LNV}} \equiv M_R \). Leptogenesis provided this information, since the baryon asymmetry is most naturally reproduced for \( M_R \geq 10^{12} \text{ GeV} \). For \( \Lambda_{\text{FV}} \) around the TeV scale, the branching ratio for \( \mu \to e\gamma \) is expected to be in the reach of the MEG experiment.

6 Conclusions

In this talk we studied leptogenesis in the MFV scenario with right-handed Majorana neutrinos degenerate in mass. Radiative corrections lift the tree-level degeneracy of right-handed neutrinos and induce mass-splittings proportional to the neutrino and charged lepton Yukawa couplings. We showed that leptogenesis is viable and most efficient at high values of right-handed neutrino...
masses ($\geq 10^{13}$ GeV), where it is driven by the mass-splittings quartic in the neutrino Yukawa couplings. As a consequence, the predictions for $\mu \rightarrow e\gamma$ are enhanced and should be observable in next experiments, at least for natural values of the scale of new physics. High energy CP-violating parameters, that disappear in the see-saw relation but take part into leptogenesis, have a significative impact on low-energy processes. The expectation $BR(\mu \rightarrow e\gamma)/BR(\tau \rightarrow \mu\gamma) \ll 1$, valid in the CP limit, is recovered in the regime of very heavy right-handed neutrinos ($M_R \gg 10^{12}$ GeV).

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1. R. S. Chivukula and H. Georgi,
SOFT LEPTOGENESIS IN THE INVERSE SEESAW

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We consider "soft leptogenesis" in the context of the inverse seesaw mechanism. In this model there are lepton number ($L$) conserving and $L$-violating soft supersymmetry-breaking $B$-terms involving the singlet sneutrinos which, together with the — generically small— $L$-violating parameter responsible of the neutrino mass, give a small mass splitting between the four singlet sneutrino states of a single generation. In combination with the trilinear soft supersymmetry breaking terms they also provide the CP violation needed to generate a lepton asymmetry in the singlet sneutrino decays. We obtain that the lepton asymmetry is proportional to the $L$-conserving soft supersymmetry-breaking $B$-term, and it is not suppressed by the $L$-violating parameters. As in the standard seesaw case, we find that this mechanism can lead to successful leptogenesis only for relatively small value soft bilinear $B$-term. The right-handed neutrino masses can be sufficiently low to elude the gravitino problem. Also the corresponding Yukawa couplings involving the lightest of the right-handed neutrinos are constrained to be $\sum |Y_{13}|^2 \lesssim 10^{-7}$, which generically implies that the neutrino mass spectrum has to be strongly hierarchical.

1 Introduction

The discovery of neutrino oscillations makes leptogenesis a very attractive solution to the baryon asymmetry problem. In the standard framework the tiny neutrino masses are generated via the (type I) seesaw mechanism and thus the new singlet neutral leptons with heavy Majorana masses can produce dynamically a lepton asymmetry through out of equilibrium decay. Eventually, this lepton asymmetry is partially converted into a baryon asymmetry due to fast $B - L$ violating sphaleron processes.

For a hierarchical spectrum of right-handed neutrinos, successful leptogenesis requires generically quite heavy singlet neutrino masses, of order $M > 2.4(0.4) \times 10^{9}$ GeV for vanishing (thermal) initial neutrino densities. The stability of the hierarchy between this new scale and the electroweak one is natural in low-energy supersymmetry, but in the supersymmetric seesaw scenario there is some conflict between the gravitino bound on the reheat temperature and the thermal production of right-handed neutrinos.

Once supersymmetry has been introduced, leptogenesis is induced also in singlet sneutrino decays. If supersymmetry is not broken, the order of magnitude of the asymmetry and the basic mechanism are the same as in the non-supersymmetric case. However, supersymmetry-breaking terms can play an important role in the lepton asymmetry generated in sneutrino decays because they induce effects which are essentially different from the neutrino ones. In brief, soft supersymmetry-breaking terms involving the singlet sneutrinos remove the mass degeneracy between the two real sneutrino states of a single neutrino generation, and provide new sources of lepton number and CP violation. As a consequence, the mixing between the
two sneutrino states generates a CP asymmetry in the decay, which can be sizable for a certain range of parameters. In particular the asymmetry is large for a right-handed neutrino mass scale relatively low, in the range $10^5 - 10^8$ GeV, well below the reheat temperature limits, what solves the cosmological gravitino problem. Moreover, contrary to the traditional leptogenesis scenario, where at least two generations of right-handed neutrinos are required to generate a CP asymmetry in (s)neutrino decays, in this new mechanism for leptogenesis the CP asymmetry in sneutrino decays is present even if a single generation is considered. This scenario has been termed “soft leptogenesis”, since the soft terms and not flavour physics provide the necessary mass splitting and CP-violating phase.

We have studied\textsuperscript{6} soft leptogenesis in the framework of an alternative mechanism to generate small neutrino masses, namely the inverse seesaw scheme\textsuperscript{7}. This scheme is characterized by a small lepton number violating Majorana mass term $\mu$, while the effective light neutrino mass is $m_{\nu} \propto \mu$. Small values of $\mu$ are technically natural, given that when $\mu \to 0$ a larger symmetry is realized: lepton number is conserved and neutrinos become massless. In the inverse seesaw scheme lepton flavour and CP violation can arise even in the limit where lepton number is strictly conserved, due to the mixing of the SU(2) doublet neutrinos with new SU(2) × U(1) singlet leptons.

As opposite to the standard seesaw case, these singlet leptons do not need to be very heavy\textsuperscript{8}, and, as a result, lepton flavour and CP violating processes are highly enhanced. In Ref.\textsuperscript{9} it was studied the possibility that the baryon asymmetry is generated in this type of models during the electroweak phase transition, in the limit $\mu = 0$. A suppression was found due to the experimental constraints on the mixing angles of the neutrinos. Therefore we considered the supersymmetric version of the model and the soft leptogenesis mechanism, since (i) in this case we expect that a CP asymmetry will be generated in sneutrino decays even with a single-generation and no suppression due to the mixing angles is expected, and (ii) this scheme provides a more natural framework for the relatively low right-handed neutrino mass scale.

2 Inverse Seesaw Mechanism

In this type of models\textsuperscript{7}, the lepton sector of the Standard Model is extended with two electroweak singlet two-component leptons per generation, i.e.,

$$L^i = \left(\nu^i_L, e^i_L, \nu^i_R, s^i_L\right)$$

We assign lepton number $L = 1$ to the singlets $s^i_L$ and $\nu^i_R$. The $(9 \times 9)$ mass matrix of the neutral lepton sector in the $\nu_L, \nu_R, s_L$ basis is given by

$$M = \begin{pmatrix} 0 & m_D & 0 \\ m_D^T & 0 & M^T \\ 0 & M & \mu \end{pmatrix}$$

where $m_D, M$ are arbitrary $3 \times 3$ complex matrices in flavour space and $\mu$ is complex symmetric. The matrix $M$ can be diagonalized by a unitary transformation, leading to nine mass eigenstates $n_\nu$: three of them correspond to the observed light neutrinos, while the other three pairs of two component leptons combine to form three quasi-Dirac leptons.

In this “inverse seesaw” scheme, assuming $m_D, \mu \ll M$ the effective Majorana mass matrix for the light neutrinos is approximately given by

$$m_{\nu} = m_D^T M^{-1} \mu M^{-1} m_D,$$

while the three pairs of heavy neutrinos have masses of order $M$, and the admixture among singlet and doublet SU(2) states is suppressed by $m_D/M$. Although $M$ is a large mass scale
suppressing the light neutrino masses, in contrast to the Majorana mass ($\Delta L = 2$) of the right-handed neutrinos in the standard seesaw mechanism, it is a Dirac mass ($\Delta L = 0$), and it can be much smaller, since the suppression in Eq. 3 is quadratic and moreover light neutrino masses are further suppressed by the small parameter $\mu$ which characterizes the lepton number violation scale.

We consider the supersymmetric version of this model. The above neutral lepton mass matrix (Eq. 2) is described by the following superpotential:

$$W = Y_{ij} N_i L_j H + \frac{1}{2} \mu_{ij} S_i S_j + M_{ij} S_i N_j ,$$

where $i, j = 1, 2, 3$ are flavour indices, $H, L_i, N_i, S_i$ are the superfields corresponding to the $SU(2)$ up-Higgs and lepton doublets, and $\nu_R^C$ and $s_L^i$ singlets respectively, and $Y_{ij}$ denote the neutrino Yukawa couplings. After spontaneous electroweak symmetry breaking, the neutrino Dirac masses are given by $(m_D)_{ij} = Y_{ij}(H)$.

The relevant soft supersymmetry breaking terms are the bilinear and trilinear scalar couplings involving the singlet sneutrino fields. We consider a simplified one generation model because a single generation of singlet sneutrinos is sufficient to generate the CP asymmetry.

$$-L_{\text{soft}} = A Y_{11} \bar{L}_1 \tilde{N} H + \tilde{m}_S^2 \tilde{S} \tilde{S}^\dagger + \tilde{m}_N^2 \tilde{N} \tilde{N}^\dagger + \tilde{m}_{SN}^2 \tilde{S} \tilde{N}^\dagger + B_S \tilde{S} \tilde{S} + B_{SN} \tilde{S} \tilde{N} + h.c. \quad (5)$$

With our lepton number assignments, the soft SUSY breaking terms which violate $L$ are $\tilde{m}_{SN}$ and $B_S$. The lagrangian describing sneutrino interactions has three independent physical CP violating phases. Furthermore, if we assume conservative values of the soft breaking terms:

$$A \sim O(m_{\text{SUSY}})$$
$$\tilde{m}_N \sim \tilde{m}_S \sim \tilde{m}_{SN} \sim O(m_{\text{SUSY}})$$
$$B_S \sim O(m_{\text{SUSY}} \mu)$$
$$B_{SN} \sim O(m_{\text{SUSY}} m) \quad (6)$$

with both, $\mu, m_{\text{SUSY}} \ll M$, we see that $B_S, \tilde{m}_N^2, \tilde{m}_S^2, \tilde{m}_{SN}^2 \ll B_{SN}$ and $\overline{M}_{SN}^2 \sim \mu M^*$. Neglecting these small soft terms, there is still one physical CP violating phase,

$$\phi = \text{arg}(AB_{SN}^* M) . \quad (7)$$

In this limit the mass degeneracy among the four sneutrino states is removed by both the $L$-violating mass $\mu$ and $L$-conserving supersymmetry breaking term $B_{SN}$. Together with the trilinear $A$ term they also provide a source of CP violation, and the mixing among the four sneutrino states leads to a CP asymmetry in their decay.

3 The CP Asymmetry

Here we present the CP asymmetry in the singlet sneutrino decays. We assume that the sneutrinos are in a thermal bath with a thermalization time $\Gamma^{-1}$ shorter than the typical oscillation times, $\Delta M_{ij}^{-1}$, therefore coherence is lost and it is appropriate to compute the CP asymmetry in terms of the mass eigenstates.

We define a fermionic and a scalar CP asymmetry in the decay of each $\tilde{N}_i$ in terms of decay
amplitudes as:

\[
\epsilon_{s_t} = \frac{\sum_k |\hat{A}_t(\tilde{N}_t \rightarrow \tilde{L}_k H)|^2 - |\hat{A}_t(\tilde{N}_t \rightarrow \tilde{L}_k^\dagger H^\dagger)|^2}{\sum_k |\hat{A}_t(\tilde{N}_t \rightarrow \tilde{L}_k H)|^2 + |\hat{A}_t(\tilde{N}_t \rightarrow \tilde{L}_k^\dagger H^\dagger)|^2}
\]

(8)

\[
\epsilon_{f_t} = \frac{\sum_k |\hat{A}_t(\tilde{N}_t \rightarrow L_k h)|^2 - |\hat{A}_t(\tilde{N}_t \rightarrow L_k^\dagger h)|^2}{\sum_k |\hat{A}_t(\tilde{N}_t \rightarrow L_k h)|^2 + |\hat{A}_t(\tilde{N}_t \rightarrow L_k^\dagger h)|^2}.
\]

(9)

To compute \(\epsilon_{s_t}\) and \(\epsilon_{f_t}\) we follow the effective field theory approach described in \(^{10}\), which considers the CP violation due to mixing of nearly degenerate states by using resummed propagators for unstable (mass eigenstate) particles. We neglect thermal corrections to the CP asymmetry from loops. We obtain the following fermionic and scalar CP asymmetries at \(T = 0\):

\[
\epsilon_{s_t} = -\epsilon_{f_t} = \hat{\epsilon}_t = -\frac{4|BSN A|}{4|M|^2 |\Gamma|^2} \sin \phi.
\]

(10)

As long as we neglect the zero temperature lepton and slepton masses and small Yukawa couplings, the phase-space factors of the final states are flavour independent.

The total asymmetry \(\epsilon_t\) generated in the decay of the singlet sneutrino \(\tilde{N}_t\) can be written in the approximate decay at rest as:

\[
\epsilon_t = \frac{\epsilon_{s_t} c_s + \epsilon_{f_t} c_f}{c_s + c_f} = \hat{\epsilon}_t \left[ \frac{(1 + n_B)^2}{(1 + n_E)^2} \right]
\]

(11)

where \(c_s, c_f\) are the phase-space factors of the scalar and fermionic channels respectively, and \(n_{B(F)}\) is the Bose(Fermi) distribution.

We find that this leptogenesis scenario presents many features analogous to soft leptogenesis in seesaw models \(^{4,5}\): (i) The CP asymmetry (Eq. 11) vanishes if \(c_s = c_f\), because then there is an exact cancellation between the asymmetry in the fermionic and bosonic channels. Finite temperature effects break supersymmetry and make the fermion and boson phase-spaces different \(c_s \neq c_f\), mainly because of the final state Fermi blocking and Bose stimulation factors. (ii) It also displays a resonance behaviour with the maximum placed at \(2|BSN|/M \sim \Gamma\). (iii) The CP asymmetry is due to the presence of supersymmetry breaking and irremovable CP violating phases, thus it is proportional to \(|BSN A| \sin \phi\).

We obtain that the CP asymmetry is not suppressed by the lepton number violating scale \(\mu\). This may seem counterintuitive. However if \(\mu = 0\) the four sneutrino states are pair degenerate, and we can choose a lepton number conserving mass basis, made of the \((L = 1)\) states

\[
\tilde{N}_1^t = \frac{1}{\sqrt{2}} \left( \tilde{S}^t - \tilde{N} \right)
\]

\[
\tilde{N}_2^t = \frac{1}{\sqrt{2}} \left( \tilde{S}^t + \tilde{N} \right)
\]

(12)

and their hermitian conjugates, with \(L = -1, \tilde{N}_1^t, \tilde{N}_2^t\). Although there is a CP asymmetry in the decay of these sneutrinos, it is not a lepton number asymmetry (since in the limit \(\mu = 0\) total lepton number is conserved) but just a redistribution of the lepton number stored in heavy sneutrinos and light lepton and slepton \(SU(2)\) doublets. At very low temperatures, \(T \ll M\), when no heavy sneutrinos remain in the thermal bath, all lepton number is in the light species and obviously if we started in a symmetric Universe with no lepton number asymmetry it will not be generated.
4 Results

Finally we quantified the conditions on the parameters which can be responsible for a successful leptogenesis: WMAP measurements in the $\Lambda$CDM model imply\textsuperscript{11} that $n_{\Sigma}/s = (8.7 \pm 0.3) \times 10^{-11}$. Assuming thermal initial sneutrino distribution and the weak wash-out regime, the final value of the baryon asymmetry generated in sneutrino decays can be parameterized as:

$$\frac{n_{\Sigma}}{s} = \frac{8}{23} \frac{n_{\Sigma}^{eq}}{s} \kappa \sum_{i} \epsilon_{i}(T_{d}),$$  

(13)

where $8/23$ is the fraction of $B - L$ asymmetry converted into baryon asymmetry by sphaleron processes, $n_{\Sigma}^{eq}/s$ is the equilibrium number density of sneutrinos normalized to the entropy density, $\epsilon_{i}(T)$ is given in Eq. 11 and $T_{d}$ is the temperature at the time of decay ($\Gamma = H(T_{d})$). Finally, $\kappa \lesssim 1$ is a dilution factor which takes into account the possible inefficiency in the production of the singlet sneutrinos, the erasure of the generated asymmetry by $L$-violating scattering processes and the temperature dependence of the CP asymmetry $\epsilon_{i}(T)$. The precise value of $\kappa$ can only be obtained from numerical solution of the Boltzmann equations. In what follows we will use an approximate constant value $\kappa = 0.2$.

Some constraints arise from the timing of the decay. First, successful leptogenesis requires the singlet sneutrinos to decay out of equilibrium: its decay width must be smaller than the expansion rate of the Universe $\Gamma < H |_{T=M}$. Second, in order for the generated lepton asymmetry to be converted into a baryon asymmetry via the $B-L$ violating sphaleron processes, the singlet sneutrino decay should occur before the electroweak phase transition ($\Gamma > H (T \sim 100 \text{ GeV})$). These two conditions determine a range for the possible values of $\sum |Y_{1k}|^2$ for a given $M$:

$$2.6 \times 10^{-21} \left(\frac{10^8 \text{ GeV}}{M}\right) < \sum |Y_{1k}|^2 < 5 \times 10^{-9} \left(\frac{M}{10^8 \text{ GeV}}\right)$$  

(14)

Eq.14 implies that the contribution of the lightest pseudo-Dirac singlet neutrino generation to the neutrino mass is negligible. Consequently, to reproduce the observed mass differences $\Delta m_{\odot}^2$ and $\Delta m_{\text{atm}}^2$, the dominant contribution to the neutrino masses must arise from the exchange of the heavier singlet neutrino states.

In Fig.1 we plot the range of parameters $\sum |Y_{1k}|^2$ and $B_{SN}$ for which enough asymmetry is generated. We show the ranges for three values of $M$ and for a generic value of $A = m_{\text{SU3Y}} = 10^9 \text{ GeV}$ and $\sin \phi = 1$. We see that this mechanism works for relatively small values of $M$ ($< 10^9 \text{ GeV}$). The smaller is $M$, the smaller are the yukawas $\sum |Y_{1k}|^2$. Also, in total analogy with the standard seesaw\textsuperscript{4,5}, the value of the soft susy-breaking bilinear $B_{SN}$, is well bellow the expected value $M m_{SU3Y}$. The reason is that, in order to generate an asymmetry large enough $B_{SN} \sim M \Gamma$, but $\Gamma$ is very small if the sneutrinos decay out of equilibrium, $\Gamma \ll 1 \text{ GeV} \left(\frac{M}{10^9 \text{ GeV}}\right)^2$.

Given the small required values of $B_{SN}$ one can question the validity of our results. In order to verify their stability we checked that as long as $|B_{SN}| \gg |B_{S}|, |\mu|^2$ the total CP asymmetry is always proportional to $B_{SN}$, and presents the same resonant behaviour, so that it is still significant only for $B_{SN} \ll M m_{SU3Y}$.

In summary in this work\textsuperscript{6} we studied the conditions for successful soft leptogenesis in the context of the supersymmetric inverse seesaw mechanism. This scheme is characterized by a small lepton number violating Majorana mass term $\mu$ with the effective light neutrino mass being $m_{\nu} \propto \mu$.

The relevant lepton asymmetry is proportional to the $L$-conserving term $B_{SN}$ and is not suppressed by any $L$-violating parameter. As in the standard see-saw case, the asymmetry displays a resonant behaviour, and can lead to successful leptogenesis only for relatively small values of $B_{SN}$. The right-handed neutrino masses are low enough to elude the gravitino problem.
Figure 1: $\sum \left| Y_{ij} \right|^2 - B_{SN}$ regions in which enough $CP$ asymmetry can be generated, and the non-equilibrium decay and decay before the electroweak phase transition conditions are verified. We take $A = m_{\tilde{SU}2} = 10^3$ GeV, $\sin \phi = 1$. The regions correspond to $M = 10^6$, $5 \times 10^7$, and $10^8$ GeV.

Also, the out of equilibrium decay condition implies that the Yukawa couplings involving the lightest of the right-handed neutrinos are constrained to be very small. This means that the contribution of the lightest pseudo-Dirac singlet neutrino generation to the neutrino mass is negligible, so the dominant contribution to the neutrino masses must arise from the exchange of the heavier singlet neutrino states. Generically this leads to a neutrino mass spectrum strongly hierarchical.

Acknowledgments

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References

VI - Neutrino Physics
CERN NEUTRINOS TO GRAN SASSO (CNGS): STATUS AND FUTURE PROTON BEAM OPTIONS

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The CERN Neutrinos to Gran Sasso project (CNGS) aims at directly detecting $\nu_\mu - \nu_\tau$ neutrino oscillations. An intense muon-neutrino beam ($10^{17} \nu_\mu / \text{day}$) is generated at CERN and directed towards the Gran Sasso National Laboratory, LNGS, in Italy, 732 km away from CERN where large and complex detectors will allow detecting, in particular, the rare tau-neutrinos. In summer 2006 the first CNGS physics runs were carried out after a successful commissioning of the CNGS facility. In the presently approved physics programme, it is foreseen to run the facility for five years with $4.5 \cdot 10^{19}$ protons/year at 400 GeV/c on the CNGS target. The maximum proton intensity to CNGS is summarized for upgrades with different beam injector scenarios.

Introduction
An overview of the CNGS neutrino beam facility\textsuperscript{1)} at CERN is shown in Figure 1. During a nominal CNGS cycle, i.e. every 6s, two SPS extractions (10.5\,\mu s each, separated by 50\,ms) of $2.4 \cdot 10^{13}$ protons each at 400\,GeV/c are sent down the 800m long proton beam line to the target. The CNGS beam is extracted from the SPS using the same extraction channel as for one of the two LHC beams. After about 100m from the extraction point, a string of switch magnets is used in order to direct the beam either to the LHC or to the CNGS target. In the CNGS graphite target, pions and kaons are produced. The positively charged $\pi/K$ are energy-selected and guided with two focusing lenses, i.e. horn and
reflector, in the direction towards Gran Sasso. In the 1000m long decay vacuum tube these particles decay into muon-neutrinos and muons. The analysis of the two muon detector stations is used to derive the intensity of the neutrino beam produced, the beam profile and an indication on the neutrino beam spectrum.

![Diagram of CNGS layout](image)

**Figure 1: CNGS layout.**

**Beam Line Commissioning**

During three weeks in July and August 2006 the primary beam line (from the SPS extraction to the target) and the secondary beam line (from the target to the muon monitors) were commissioned. The beam monitoring equipment was tested and measurements were compared with simulations.

The proton beam position monitors (BPMs) revealed that the primary beam line was well tuned over its 800m after only a minor magnetic correction.

The maximum trajectory beam excursion is well within the specification. The beam position stability onto the target has been averaged over several days and measured to be $\sim 50 \mu$m r.m.s.² ³ ⁴.

The muon detectors provide on-line feedback for the quality control of the neutrino beam. The muon detectors are nitrogen-filled, sealed ionization chambers and are designed to measure up to $10^8$ muons per cm² and per extraction. There are 41 fixed detectors and one movable detector installed in each of the two CNGS muon
detector chambers. They are assembled in a cross-shaped array to provide the muon intensity and the vertical and horizontal muon profile.

Using the muon monitor profile information (centroid and symmetry), we are able to optimize the alignment between the proton beam, the target and the horn/reflector (see Fig. 2). The muon measurements confirmed the stability of the proton beam and showed that the beam equipment downstream of the target performed as expected\textsuperscript{6}.

\textbf{Figure 2:} Vertical muon profile in the second muon detector chamber for different beam versus target alignments (0mm, +1mm, -1.5mm). Each point corresponds to a detector and shows the measured induced charges per proton on target (\textasciitilde2700 charges/muon are produced in a detector). The statistical accuracy is limited (as indicated by the fluctuations of the data points) because the dynamic range of the readout electronics is adjusted to the nominal CNGS intensity ($2.4 \times 10^{13}$ p.o.t./extraction). The measurements were performed with $\approx 2 \times 10^{13}$ p.o.t./extraction.
CNGS Operation

The first CNGS physics run started as scheduled on 18\textsuperscript{th} August 2006 and lasted 12 days. Another run dedicated for beam optimizing has been foreseen for 2 weeks starting on 26\textsuperscript{th} October 2006. This run had to be stopped after two days due to a water leak in the water cooling circuit of the reflector. During the physics run in total $\sim8.5\cdot10^{17}$ protons were delivered to the target. The average beam intensity during the physics run was at the order of $\sim1.4\cdot10^{13}$ protons on target (p.o.t.) per extraction at 400 GeV/c. The maximum beam intensity reached in 2006 was $1.75\cdot10^{13}$ p.o.t. per extraction.

Analysis of the Maximum Potential Proton Flux to CNGS

The limitations of the proton flux which can be sent to the CNGS facility have been studied together with the estimation of the maximum attainable flux\textsuperscript{7} for two scenarios: (a) operating with the present injector chain, and (b) operating with new injectors - LINAC4, SPL and PS2, as proposed by the PAF working group\textsuperscript{8}. The proton flux was calculated with the assumption of 200 days of operation per year and 80\% injector availability. The intensity limitations coming from equipment (as designed) in the CNGS facility are summarized in Table1.

With the present injectors, the maximum achievable beam intensity in the SPS is estimated at $5.7\cdot10^{13}$ protons per cycle with machine improvements aimed mainly at beam loss reduction. With $5.7\cdot10^{13}$ protons per cycle, about $1.1\cdot10^{20}$ p.o.t. per year can be obtained in the scenario of 85\% of the SPS beam time dedicated to CNGS.

With new injectors (after 2016), assuming a major upgrade of the SPS RF power capability and solutions to the e-cloud problems, beam instabilities and the heating by the beam current of numerous equipment in the SPS, the maximum number of protons accelerated per SPS cycle can reach $1\cdot10^{14}$. The integrated proton flux can potentially attain $2.4\cdot10^{20}$ p.o.t. per year if the SPS is dedicated 85\% of the time to CNGS and the SPS cycle length is 4.8 s. Beyond $7\cdot10^{13}$ p.o.t. per cycle and/or
beyond $1.38 \times 10^{20}$ p.o.t. per year, the design limit of most of the secondary beam line equipment is reached and a full redesign is mandatory to address higher intensities.

<table>
<thead>
<tr>
<th>Intensity limitation</th>
<th>Protons per batch</th>
<th>Protons per 6s cycle</th>
<th>Proton flux [protons on target per year]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Radiation protection calculation and optimisation</td>
<td>$3.5 \times 10^{13}$</td>
<td>$7 \times 10^{13}$</td>
<td>Soil/concrete activation: $4.5 \times 10^{19}$</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>Air/water activation: $7.6 \times 10^{19}$</td>
</tr>
<tr>
<td>Target design</td>
<td>$1.4 \times 10^{14}$</td>
<td></td>
<td>Radiation damage: $2 \times 10^{20}$</td>
</tr>
<tr>
<td>Horn design</td>
<td>$3.5 \times 10^{13}$</td>
<td>$1 \times 10^{14}$</td>
<td>Horn air cooling system: $1.38 \times 10^{20}$</td>
</tr>
<tr>
<td>Shielding, Decay Tube, Hadron stop design</td>
<td></td>
<td></td>
<td>Cooling: $1.38 \times 10^{20}$</td>
</tr>
</tbody>
</table>

Operation beyond the present nominal CNGS parameters implies that all radiation protection calculations/studies have to be revised, including radioactive waste studies and the corresponding area classifications. A new INB approval from IRSN will be mandatory for operation beyond nominal parameters. The replacement of the initial equipment in the target cavern after the nominal 5 year run for CNGS will be extremely challenging and its many years radiation down-cooling time must be estimated. In the second scenario (new injectors) the CNGS facility itself will need a major rebuild, because of the difficulty to replace the first generation of equipment in the target building (activation and risk of contamination) and also because of the
need to reassess all radiation protection issues and to dimension the new equipment, tunnel and cavern accordingly.

Summary
The CNGS facility has been successfully commissioned and operated for a first physics run below nominal intensity. The true challenge of CNGS starts now, with the planned continuous high intensity operation, resulting in high radiation levels in much of the CNGS area and in fatigue effects on target, horn and reflector from pulsed operation and from the beam impact on the equipment.

The experience that will be gained operating the present CNGS facility will be crucial to benchmark the design values, confirm or improve the models used in simulations, and hence contribute to the design of any upgraded future facility.

Acknowledgment
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OPERA first events from the CNGS neutrino beam

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The aim of the OPERA experiment is to search for the appearance of the tau neutrino in the quasi pure muon neutrino beam produced at CERN (CNGS). The detector, installed in the Gran Sasso underground laboratory 730 km away from CERN, consists of a lead/emulsion target complemented with electronic detectors. A report is given on the detector status (construction, data taking and analysis) and on the first successful 2006 neutrino runs.

Introduction

In the last decades solar and atmospheric neutrino experiments observed deficits in the measured fluxes which are all well reproduced in a neutrino oscillations model, implying non vanishing, not degenerate neutrino masses and neutrino mixing. Within such hypothesis weak interactions eigenstates differ from the mass eigenstates. The mixing can be parametrized in an unitary matrix whose parameters (3 angles and 1 or 3 phases depending on the Dirac or Majorana nature of neutrinos) associated to the square masses differences $\Delta m^2$ drive the amplitude of the disappearance ($P(\nu_\alpha \rightarrow \nu_\beta)$) or survival ($P(\nu_\alpha \rightarrow \nu_\alpha)$) probabilities. The major experimental results for solar neutrinos come from radiochemic experiments$^{1,2,3}$ or large Čerenkov detectors$^{4,5,6,7}$. For atmospheric neutrinos they come from Čerenkov and calorimeters$^{8,9}$. These results were confirmed with reactor$^{10,11}$ and long baseline experiments$^{12,13}$. All these experiments are however of “disappearance” type, they compare the measured flux at a far distance with either one at a close position (in the case of long baseline experiments) or with the predicted one (in the reactor, “solar” and “atmospheric” experiments).

The OPERA experiment$^{14}$ has been designed to perform an unique appearance observation of the oscillation products to confirm (or infirm) the neutrino oscillation hypothesis in the atmospheric sector through the $\nu_\mu \rightarrow \nu_\tau$ channel and also to set limits on the $\theta_{13}$ angle through the $\nu_\mu \rightarrow \nu_\tau$ channel. This article reports the first observed neutrino events from the CNGS (CERN to Gran Sasso) beam by the OPERA experiment.

1 The OPERA experiment

1.1 The CNGS programme

The CNGS$^{15}$ programme of neutrino beam from CERN to Gran Sasso has been approved in 1999. OPERA was approved as “CNGS1 experiment” in 2001. From CERN to Gran Sasso the neutrinos time of flight is 2.44 ms and their average direction makes a 3° angle w.r.t. the horizontal due to the earth curvature. The main features of the beam have been presented in
details in this conference. The beam has been optimized to maximize the number of \( \tau \) events in the detector (convolution of the neutrino flux, the disappearance probability \( P(\nu_\mu \to \nu_\tau) \) and the detection efficiency). The neutrino average energy is 17 GeV. The \( \bar{\nu}_\mu \) contamination is \( \sim 4\% \), the \( \bar{\nu}_e \) (\( \bar{\nu}_\mu \)) is < 1\% and the number of \( \nu_\tau \) is negligible. The expected beam intensity is \( 4.5 \times 10^{19} \) p.o.t./year.

1.2 The detection technique

The challenge of the experiment is to measure the appearance of \( \nu_\tau \) from \( \nu_\mu \) oscillations through CC \( \tau \) interactions. The events induced by the short-lived \( \tau \) have a characteristic topology (with a "kink" due to the presence of undetected neutrinos in the \( \tau \) decay) but extends over \( \sim mm^3 \) typical volumes.

ECC technique The detector should therefore match a large mass for statistics, a high spatial resolution and high rejection power to limit background contamination. These requirements are satisfied using the proven ECC (Emulsion Cloud Chamber) technique which already worked successfully in the DONUT experiment. The passive target consists of lead plates. Particles are tracked in nuclear emulsion films with a sub-micrometric intrinsic resolution. 57 emulsion films are assembled and interspaced with 56 lead plates 1 mm wide in a detector basic cell called "brick". A brick is a 12.7 \times 10.2 cm\(^2\) object with a thickness along the beam direction of 7.5 cm (about 10 radiation lengths). Its weight is about 8.3 kg. Bricks are assembled in 31 walls (52 \times 64 bricks) separated by electronic detectors planes to trigger the event and identify the brick with the interaction vertex.

The readout sequence of OPERA events is quasi-online. Once identified the brick is extracted from the detector, emulsions are developed and scanned by automatic microscopes. Scanning performs detailed tracking, vertex location, particle identification, momentum measurement through Multiple Coulomb Scattering, decay kink search. The data are complemented with the momentum, energy and charge measurements done by the electronic detectors.

Expected performances At the expected nominal beam intensity and for five years data taking a total of 31000 charged and neutral current interactions is expected in the nominal mass target of OPERA. Among these 35 (214) CC \( \nu_\tau \) interactions are expected for oscillation parameter values \( \Delta m^2_{3\alpha}=2 (3) \times 10^{-3} \) eV\(^2\) and \( \sin^2 2\theta_{23}=1 \). The overall detection efficiencies have been estimated in Monte Carlo simulations upgraded by dedicated tests (vertex searches, \( e/\pi \) and \( \pi/\mu \) separation, large angle muon scattering...). The nominal expected number of \( \tau \) events ranges from 11 to 16 events with the same parameters sets. The physics channel considered so far are \( \tau \to e, \tau \to \mu \) and \( \tau \to h\) (1 or 3 prongs). The background is expected to be < 1. In the sub-dominant \( \nu_\mu \to \nu_e \) channel OPERA should, in the same run conditions, set a limit of \( \sin^2 2\theta_{13} < 0.06 \) (90\% C.L.) assuming \( \Delta m^2 = 2.5 \times 10^{-3} \) eV\(^2\).

1.3 The OPERA detector

The OPERA detector is divided into two Super-Modules consisting of a target section followed by a muon spectrometer (see Picture 1). The target section is made of 31 brick walls each one being followed by a highly segmented scintillator tracker plane. A large VETO plane is placed in front of the detector to further discriminate beam events from horizontal cosmics. The construction of the experiment started in Spring 2003. The two instrumented magnets were completed in May 2004 and beginning of 2005 respectively. In Spring 2006 all scintillator planes were installed.
The electronic detectors  The target tracker covers a total area of 7000 m$^2$ and is built of 32000 scintillator strips, each $7 \text{ m}$ long and of 25 mm x 15 mm cross section. Along the strip, a wavelength shifting fiber of 1 mm diameter transmits the light signals to both ends, read out by 992 multi-anode (64 channels) PMT's from Hamamatsu. A 32 channel front-end electronics ASIC$^{19}$ has been developed which allows individual gain corrections (with a dynamic range of 1 \div 4) and auto-triggered readout sequence in a standard dual shaper scheme (fast and slow) and track & hold logic.

The muon spectrometer consists of a large $8 \times 8 \text{ m}^2$ dipolar magnet delivering a magnetic field of 1.55 T and instrumented with RPC's and drift tubes. Each magnet arm consists of twelve 5 cm thick iron slabs, alternating with RPC planes. This sandwich structure allows the tracking in the magnetic field to identify the muons and to determine their momentum. In addition the precision tracker$^{20}$ measures the muon track coordinates in the horizontal plane. It is made of $8 \text{ m}$ long drift tubes with an outer diameter of 38 mm. The charge misidentification is expected to be 0.1 % - 0.3 % in the relevant momentum range which is efficient enough to minimize the background originating from the charm particles produced in $\nu_p$ interactions. With the muon spectrometer a momentum resolution of $\Delta p/p \leq 0.25$ for all muon momenta $p$ up to a maximum of $p = 25 \text{ GeV/c}$ can be achieved.

"Bricks" production  In total $\sim 200000$ bricks should be nominally produced and installed in OPERA. The production is performed by a dedicated apparatus called Brick Assembly Machine (BAM) installed in Gran Sasso. It is a chain of different stations (for piling, pressing, wrapping etc) using robots operating in dark rooms with controlled environment. The production is around $\sim 2 \text{ bricks per minute}$. Once bricks are produced they are placed in dedicated mechanical structures called drums (9 rows of 26 bricks) and inserted from there to the detector by the Brick Manipulating System (BMS). The BMS has two equivalent structures (one per side) consisting of a brick storage place (carousel) where the drums are exchanged and a moving robot along the side of the experiment. The mobile part of the BMS can reach the desired row and plane with a sub-millimetric accuracy. The robot has a mobile bridge on which the bricks are pushed inside the detector by a pushing arm to their desired position. The brick extraction is performed by a vacuum sucker located on the front of a small vehicle.
2 Data taking and analysis

DAQ and on-line analysis  The OPERA DAQ system is based on a so-called “smart” sensor concept on an Ethernet network. The principle is to implement a local micro-processor as close as possible to the front-end electronics and to access it for configuration and/or data transmission through Ethernet. The core of this architecture is a small processor board which hosts a sequencer (FPGA of the Altera cyclone family), a micro-processor (32 bits RISC Etrax100lx processor from Axis) and an intermediate FIFO. The FPGA manages the full readout sequence, the data timestamping with a 10 ns accuracy, the data formatting and pre-processing (pedestals and zero suppression, local histogramming) and the data transfer to the intermediate buffer. The processor runs “sensor” applications communicating with “daq” servers and developed within the CORBA framework implemented in C++.

Each sensor is plugged to a sub-detector specific motherboard and is seen as a standard Ethernet node over a large network. In total 1153 sensors are connected for a total of 100000 readout channels. The synchronization of each individual clock is performed through a specific bi-directional bus starting from a GPS PCI board developed on purpose. A common 20 MHz clock embedding specific signals is send to the sensors with a measurement of time propagation delays for off-line correction.

Data analysis and reconstruction is performed continuously on-line within the Opera software framework based on ROOT.

Off-line emulsion scanning  After development emulsions are scanned by automatic microscopes whose nominal speed is higher than $\sim 20$ cm$^2$/h per emulsion layer (44 $\mu$m thick). There are two different approaches developed by the OPERA collaboration, in Europe (ESS$^{21}$, based on software image reconstruction) and in Japan (S-UTS$^{22}$ based on hard-coded algorithms). Next picture displays an example of both systems. The scanning sequence proceeds with the division of the emulsion thickness into $\sim 16$ tomographic images by focal plane adjustment, images digitization and track finding algorithms. Track grains are identified and separated from “fog” grains and associated into “micro-tracks”. Examples of reconstructed tracks during 2006 runs will be given below.

![Picture](image1.png)

Figure 2: Pictures of one of the ESS microscopes (left) and of the S-UTS (right).

3 First CNGS neutrino events

Beam structure reconstruction  During the first CNGS run in August 2006 319 neutrino events were collected with an estimated systematic error of 5% (see Fig.3 left for a typical display). The events were selected by a comparison of their absolute timestamps w.r.t. the beam time information in quite large coincidence window (1ms). The beam spill time structure reconstructed in OPERA is displayed in Fig.3 (right).
Cosmics vs neutrino events  Beam events have an average direction close to the horizontal one (3.3° angle) whereas cosmics have large angles distributions. This is shown in the distributions of Fig. 4. A gaussian fit to the central distribution leads to a mean angle of 3.4±0.3° in agreement with the expected value.

Emulsions matching  Some clean tracks have been followed in emulsion sheets inside the detector. A typical display of the hits reconstructed in the emulsions and in the TT is given in Fig. 5.
Conclusions and perspectives

OPERA performed the first detection of neutrino events from the long baseline CERN CNGS beam in the underground Gran Sasso laboratory. 319 neutrino-induced events were collected for an integrated intensity of $7.6 \times 10^{17}$ p.o.t. in agreement with the expectations. The reconstructed zenith-angle distributions and the time structure of the events demonstrate the capability of the electronic detectors, build up during the last three years, to reach the experiment goals. The association of tracks between electronic detectors and emulsion sheets has been also successfully performed. The collaboration is facing the last large effort of brick production and insertion and is preparing next neutrinos runs in fall 2007 for physics commissioning.

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RESULTS FROM HARP
AND THEIR IMPLICATIONS FOR NEUTRINO PHYSICS

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Recent results from the HARP experiment on the measurements of the double-differential production cross-section of pions in proton interactions with beryllium, carbon and tantalum targets are presented. These results are relevant for a detailed understanding of neutrino flux in accelerator neutrino experiments MiniBooNE/SciBooNE, for a better prediction of atmospheric neutrino fluxes as well as for an optimization of a future neutrino factory design.

1 The HARP experiment

The HARP experiment \(^{1,2}\) at the CERN PS was designed to make measurements of hadron yields from a large range of nuclear targets and for incident particle momenta from 1.5 GeV/c to 15 GeV/c. The main motivations are the measurement of pion yields for a quantitative design of the proton driver of a future neutrino factory, a substantial improvement in the calculation of the atmospheric neutrino flux and the measurement of particle yields as input for the flux calculation of accelerator neutrino experiments, such as K2K\(^{3,4}\), MiniBooNE\(^{5}\) and SciBooNE\(^{6}\).

The HARP experiment makes use of a large-acceptance spectrometer consisting of a forward and large-angle detection system. A detailed description of the experimental apparatus can be found in Ref. \(^{2}\). The forward spectrometer — based on large area drift chambers \(^{7}\) and a dipole magnet complemented by a set of detectors for particle identification (PID): a time-of-flight wall \(^{8}\) (TOFW), a large Cherenkov detector (CHE) and an electromagnetic calorimeter — covers polar angles up to 250 mrad which is well matched to the angular range of interest for the measurement of hadron production to calculate the properties of conventional neutrino beams. The large-angle spectrometer — based on a Time Projection Chamber (TPC) located inside a solenoidal magnet — has a large acceptance in the momentum and angular range for the pions relevant to the production of the muons in a neutrino factory. It covers the large majority of the pions accepted in the focusing system of a typical design. The neutrino beam of a neutrino factory originates from the decay of muons which are in turn the decay products of pions.
2 Results obtained with the HARP forward spectrometer

The first HARP physics publication \(^9\) reported measurements of the \(\pi^+\) production cross-section from an aluminum target at 12.9 GeV/c proton momentum. This corresponds to the energies of the KEK PS and the target material used by the K2K experiment. The results obtained in Ref. \(^9\) were subsequently applied to the final neutrino oscillation analysis of K2K \(^4\), allowing a significant reduction of the dominant systematic error associated with the calculation of the so-called far-to-near ratio (see \(^9\) and \(^4\) for a detailed discussion) and thus an increased K2K sensitivity to the oscillation signal.

A detailed description of established experimental techniques for the data analysis in the HARP forward spectrometer can be found in Ref. \(^9,10\). Our next goal is to contribute to the understanding of the MiniBooNE and SciBooNE neutrino fluxes. They are both produced by the Booster Neutrino Beam at Fermilab which originates from protons accelerated to 8.9 GeV/c by the booster before being collided against a beryllium target. As was the case for the K2K beam, a fundamental input for the calculation of the resulting \(\nu\mu\) flux is the measurement of the \(\pi^+\) cross-sections from a thin 5\% nuclear interaction length (\(\lambda\)) beryllium target at 8.9 GeV/c proton momentum, which is presented here and in the forthcoming HARP publication \(^11\).

With respect to our first published physics paper \(^9\), a number of improvements to the analysis techniques and detector simulation have been made. The most important improvements introduced in this analysis compared with the one presented in Ref. \(^9\) are:

- An increase of the track reconstruction efficiency which is now constant over a much larger kinematic range and a better momentum resolution coming from improvements in the tracking algorithm;
- Better understanding of the momentum scale and resolution of the detector, based on data, which was then used to tune the simulation. This results in smaller systematic errors associated with the unsmearing corrections determined from Monte Carlo;
- New particle identification hit selection algorithms both in the TOFW and in the CHE resulting in much reduced background and negligible efficiency losses;
- Significant increases in Monte Carlo production have also reduced uncertainties from Monte Carlo statistics and allowed studies which have reduced certain systematics.

It is important to point out that an analysis incorporating these improvements yields results for the aluminum data fully consistent with those published in Ref. \(^9\).

The absolutely normalized double-differential cross-section for the process \(p + Be \rightarrow \pi^+ + X\) can be expressed in bins of pion kinematic variables in the laboratory frame, \((p_\pi, \theta_\pi)\), as

\[
\frac{d^2\sigma^{\pi^+}}{dpd\Omega}(p_\pi, \theta_\pi) = \frac{A}{N_A t} \cdot \frac{1}{\Delta p \Delta \Omega} \cdot \frac{1}{N_{pot}} \cdot N^{\pi^+}(p_\pi, \theta_\pi),
\]

where:

- \(\frac{d^2\sigma^{\pi^+}}{dpd\Omega}\) is the cross-section in \(\text{cm}^2/(\text{GeV/c})/\text{sr}\) for each \((p_\pi, \theta_\pi)\) bin covered in the analysis.
- \(\frac{1}{N_A t}\) is the reciprocal of the number density of target nuclei for Be \((1.2349 \cdot 10^{23} \text{ per cm}^3)\).
- \(t\) is the thickness of the beryllium target along the beam direction. The thickness is measured to be 2.046 cm with a maximum variation of 0.002 cm.
- \(\Delta p\) and \(\Delta \Omega\) are the bin sizes in momentum and solid angle, respectively.\(^a\)

\(^a\Delta p = p_{\text{max}} - p_{\text{min}}; \Delta \Omega = 2\pi(\cos(\theta_{\text{min}}) - \cos(\theta_{\text{max}}))\)
Table 1: Total number of events in the 8.9 GeV/c beryllium 5% $\lambda$ target and empty target data sets, and the number of protons on target as calculated from the prescaled trigger count.

<table>
<thead>
<tr>
<th>Data Set</th>
<th>8.9 GeV/c Be 5% $\lambda$</th>
<th>8.9 GeV/c Empty Target</th>
</tr>
</thead>
<tbody>
<tr>
<td>protons on target</td>
<td>13,074,880</td>
<td>1,990,400</td>
</tr>
<tr>
<td>total events processed</td>
<td>4,682,911</td>
<td>413,095</td>
</tr>
<tr>
<td>events with accepted beam proton</td>
<td>2,277,657</td>
<td>200,310</td>
</tr>
<tr>
<td>beam proton events with FTP trigger</td>
<td>1,518,683</td>
<td>91,690</td>
</tr>
<tr>
<td>total good tracks in fiducial volume</td>
<td>95,897</td>
<td>3,110</td>
</tr>
</tbody>
</table>

- $N_{\text{pot}}$ is the number of protons on target after event selection cuts.
- $N^{\pi^+}(p_\pi, \theta_\pi)$ is the yield of positive pions in bins of true momentum and angle in the laboratory frame.

Eq. 1 can be generalized to give the inclusive cross-section for a particle of type $\alpha$

$$
\frac{d^2\sigma^{\alpha}}{dpd\Omega}(p, \theta) = \frac{A}{N_{\text{pot}}} \cdot \frac{1}{\Delta p \Delta \Omega} \cdot \frac{1}{M^{-1}_{p\alpha p'\theta'\alpha'}} \cdot N^{\alpha'}(p', \theta') ,
$$

where reconstructed quantities are marked with a prime and $M^{-1}_{p\alpha p'\theta'\alpha'}$ is the inverse of a matrix which fully describes the migrations between bins of true and reconstructed quantities, namely: lab frame momentum, $p$, lab frame angle, $\theta$, and particle type, $\alpha$.

There is a background associated with beam protons interacting in materials other than the nuclear target (parts of the detector, air, etc.). These events are subtracted by using data collected without the nuclear target in place where one has been careful to normalize the sets to the same number of protons on target. This procedure is referred to as the `empty target subtraction':

$$
N^{\alpha'}(p', \theta') \rightarrow [N^{\alpha'}_{\text{target}}(p', \theta') - N^{\alpha'}_{\text{empty}}(p', \theta')] .
$$

The event selection is performed in the following way: a good event is required to have a single, well reconstructed and identified beam particle impinging on the nuclear target. A downstream trigger in the forward trigger plane (FTP) is also required to record the event, necessitating an additional set of unbiased, pre-scaled triggers for absolute normalization of the cross-section. These pre-scale triggers (1/64 for the 8.9 GeV/c Be data set) are subject to exactly the same selection criteria for a `good' beam particle as the event triggers allowing the efficiencies of the selection to cancel, thus adding no additional systematic uncertainty to the absolute normalization of the result. Secondary track selection criteria have been optimized to ensure the quality of the momentum reconstruction as well as a clean time-of-flight measurement while maintaining high reconstruction and particle identification efficiencies. The results of the event and track selection in the beryllium thin target data set are shown in Table 1.

The double-differential inelastic cross-section for the production of positive pions from collisions of 8.9 GeV/c protons with beryllium have been measured in the kinematic range from 0.75 GeV/c $\leq p_\pi \leq$ 6.5 GeV/c and 0.030 rad $\leq \theta_\pi \leq$ 0.210 rad, subdivided into 13 momentum and 6 angular bins. Systematic errors have been estimated. A full (13 $\times$ 6)$^2 = 6048$ element covariance matrix has been generated to describe the correlation among bins. The data are presented graphically as a function of momentum in 30 mrad bins in Fig. 1. To characterize the uncertainties on this measurement we show the diagonal elements of the covariance matrix plotted on the data points in Fig. 1. A typical total uncertainty of 9.8% on the double-differential cross-section values and a 4.9% uncertainty on the total integrated cross-section are obtained.
Figure 1: HARP measurements of the double-differential production cross-section of positive pions, $d^2\sigma^{\pi^+}/dpd\Omega$, from 8.9 GeV/c protons on 5% $\lambda_1$ beryllium target as a function of pion momentum, $p$, in bins of pion angle, $\theta$, in the laboratory frame. The error bars shown include statistical errors and all (diagonal) systematic errors. The dotted histograms show the Sanford-Wang parametrization that best fits the HARP data.

Table 2: Sanford-Wang parameters and errors obtained by fitting the dataset. The errors refer to the 68.27% confidence level for seven parameters ($\Delta\chi^2 = 8.18$).

<table>
<thead>
<tr>
<th>Parameter (Units)</th>
<th>$c_1$ (mb GeV/c mrad$^{-1}$)</th>
<th>$c_2$ (mb GeV/c mrad$^{-1}$)</th>
<th>$c_3$ (mb GeV/c mrad$^{-1}$)</th>
<th>$c_4$ = $c_3$</th>
<th>$c_5$ (mb GeV/c mrad$^{-1}$)</th>
<th>$c_6$ (mb GeV/c mrad$^{-1}$)</th>
<th>$c_7$ (mb GeV/c mrad$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Value</td>
<td>(61.8 ± 10.8)</td>
<td>(66.7 ± 10.2)</td>
<td>(86.6 ± 20.2)</td>
<td>(7.44 ± 0.35) x 10$^{-4}$</td>
<td>(5.09 ± 0.40)</td>
<td>(0.187 ± 0.023)</td>
<td>(41.8 ± 13.6)</td>
</tr>
</tbody>
</table>

Sanford and Wang\textsuperscript{12} have developed an empirical parametrization for describing the production cross-sections of mesons in proton-nucleus interactions. This parametrization has the functional form:

$$
\frac{d^2\sigma(p+\Lambda \rightarrow \pi^+ + X)}{dpd\Omega}(p, \theta) = \exp[c_1 - c_3 \frac{p^4}{p_{beam}^{4.4}} - c_6 (p - c_7 p_{beam} \cos^2 \theta)] p^{c_1} (1 - \frac{p}{p_{beam}}),
$$

where $X$ denotes any system of other particles in the final state, $p_{beam}$ is the proton beam momentum in GeV/c, $p$ and $\theta$ are the $\pi^+$ momentum and angle in units of GeV/c and radians, respectively, $d^2\sigma/(dpd\Omega)$ is expressed in units of mb/(GeV/c mrad), $d\Omega \equiv 2\pi d(\cos \theta)$, and the parameters $c_1, \ldots, c_7$ are obtained from fits to meson production data.

The $\pi^+$ production data reported here have been fitted to this empirical formula (Eq. 4). In the $\chi^2$ minimization, the full error matrix was used. The best-fit values of the Sanford-Wang parameters are reported in Table 2, together with their errors.

The MiniBooNE neutrino beam is produced from the decay of $\pi$ and K mesons which are produced in collisions of 8.9 GeV/c protons from the Fermilab Booster on a 71 cm beryllium target. The neutrino flux prediction is generated using a Monte Carlo simulation. In this simulation the primary meson production rates are taken from a fit of existing data with a Sanford-Wang empirical parametrization in the relevant region. The results presented here, being for protons at exactly the booster beam energy, are then a critical addition to these global fits. The kinematic region of the measurements presented here contains 80.8% of the pions contributing to the neutrino flux in the MiniBooNE detector.

A similar analysis has been performed using the HARP forward spectrometer for the measurement of the double-differential production cross-section of $\pi^+$ in the collision of 12 GeV/c...
protons with a thin 5% $\lambda_1$ carbon target. The results are shown in Fig. 2. These measurements are important for a precise calculation of the atmospheric neutrino flux and for a prediction of the development of extended air showers.

3 Results obtained with the HARP large-angle spectrometer

First results on the measurements of the double-differential cross-section for the production of charged pions in proton–tantalum collisions emitted at large angles from the incoming beam direction have been obtained recently. The pions were produced by proton beams in a momentum range from 3 GeV/c to 12 GeV/c hitting a tantalum target with a thickness of 5% $\lambda_1$. The angular and momentum range covered by the experiment (100 MeV/c $\leq p < 800$ MeV/c and $0.35 \text{ rad} \leq \theta < 2.15 \text{ rad}$) is of particular importance for the design of a neutrino factory. Track recognition, momentum determination and particle identification were all performed based on the measurements made with the TPC. Results for the double-differential cross-sections $d^2\sigma/dp d\theta$ at four incident proton beam momenta (3 GeV/c, 5 GeV/c, 8 GeV/c and 12 GeV/c) are shown in Fig. 3.

Similar analyses are being performed for the Be, C, Cu, Sn and Pb targets using the same detector, which will allow a study of A-dependence of the pion yields with a reduced systematic uncertainty to be performed.

4 Conclusions

Measurements of the double-differential production cross-section of positive pions in the collision of 8.9 GeV/c protons with a beryllium target have been presented. The data have been reported in bins of pion momentum and angle in the kinematic range from 0.75 GeV/c $\leq p_\pi \leq 6.5$ GeV/c and 0.030 rad $\leq \theta_{p_\pi} \leq 0.210$ rad. A systematic error analysis has been performed yielding an average point-to-point error of 9.8% (statistical + systematic) and an overall normalization error of 2%. The data have been fitted to the empirical parameterization of Sanford and Wang and the resulting parameters provided. These production data have direct relevance for the prediction of a $\nu_\mu$ flux for MiniBooNE and SciBooNE experiments.

Preliminary results for the measurement of the double-differential production cross-section of $\pi^\pm$ in the collision of 12 GeV/c protons with a carbon target have been presented.
Figure 3: Double-differential cross-sections for $\pi^+$ (left) and $\pi^-$ (right) production in $p$-Ta interactions as a function of momentum displayed in different angular bins (shown in mrad in the panels). The results are given for all incident beam momenta (filled triangles: 3 GeV/c, open triangles: 5 GeV/c, filled rectangles: 8 GeV/c, open circles: 12 GeV/c). The error bars take into account the correlations of the systematic uncertainties.

First results on the production of pions at large angles with respect to the beam direction for protons of 3 GeV/c, 5 GeV/c, 8 GeV/c and 12 GeV/c impinging on a thin tantalum target have been described. These data can be used to make predictions for the fluxes of pions to enable an optimized design of a future neutrino factory.

Acknowledgments

It is a pleasure to thank the organizers for the financial support which allowed me to participate in the conference and to present these results on behalf of the HARP Collaboration.

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Final results from K2K and status of T2K

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The K2K (KEK-to-Kamioka) neutrino oscillation experiment is the first accelerator-based experiment with more than 100 km baseline. Data taking of K2K started in 1999 and finished in 2004. Final results for oscillation analyses are presented. The next generation long baseline experiment in Japan, T2K (Tokai-to-Kamioka), which aims for the observation of $\nu_e$ appearance from $\nu_\mu$ beam with an order of magnitude improved sensitivity, is in preparation for startup in April 2009. The status of experiment and related efforts are reported.

1 Introduction

In early 1990s, ‘anomaly’ of $\nu_\mu/\nu_e$ ratio in atmospheric neutrino was reported by Kamiokande\textsuperscript{1} and other experiments. In 1998, Super-Kamiokande (Super-K) collaboration announced famous evidence for neutrino oscillation in atmospheric neutrinos\textsuperscript{2,3}. Together with excellent results from solar (and later reactor) neutrino experiments, neutrino oscillation rapidly became ‘standard’ explanation of former neutrino ‘anomalies.’ However, in order to firmly establish this picture, it was necessary to have independent experiments, especially with controlled, artificial neutrino source. K2K (KEK-to-Kamioka) was the first experiment which confirmed the neutrino oscillation using accelerator-produced neutrino beam.

The neutrino oscillation among ‘standard’ three flavors can be characterized by three phases ($\theta_{12}, \theta_{23}, \theta_{13}$) and difference of mass-squared between corresponding mass eigenstates ($\Delta m_{12}^2$ and $\Delta m_{23}^2$). With the approximation of two-flavor oscillation, the survival probability of muon neutrino with energy $E_\nu$ (GeV) after traveling distance $L$ (km) is written as

$$P(\nu_\mu \rightarrow \nu_\mu) = 1 - \sin^2 2\theta_{23} \sin^2 [1.27 \times \Delta m_{23}^2 \times L(\text{km})/E_\nu(\text{GeV})],$$

(1)

where $\theta_{23}$ and $\Delta m_{23}^2$ are the mixing angle and mass-squared difference between second and third generation neutrinos, respectively. For this mode, $\nu_\mu$ is considered to oscillate primarily into
\[ \nu_e \] With neutrinos below \( \tau \) production energy threshold, this is observed as ‘disappearance’ of charged current events.

One can also consider ‘appearance’ of \( \nu_e \) from \( \nu_\mu \) as a sub-dominant contribution, with probability
\[
P(\nu_\mu \rightarrow \nu_e) = \sin^2 2\theta_{23} \sin^2 \theta_{13} \sin^2 [1.27 \times \Delta m^2_{23} \times L \text{(km)}/E_\nu \text{(GeV)}].
\]
\[ \text{(2)} \]

This mode is sensitive to \( \theta_{13} \), the last mixing angle to be measured and primary goals of next generation neutrino oscillation experiments.

In this article, we report final results from K2K on both \( \nu_\mu \) disappearance and \( \nu_e \) appearance analyses, as well as neutrino-nucleus interaction studies. Then, the status and prospects of the next long baseline neutrino oscillation experiment in Japan, T2K (Tokai-to-Kamioka), will be given.

2 K2K – first long baseline experiment


2.1 Experimental setup

A schematic picture of the K2K beamline is shown in Fig 1. Almost (\( \geq 98\% \)) pure \( \nu_\mu \) beam

![Figure 1: Schematic layout of K2K experiment.](image)

with a mean energy of 1.3 GeV is produced by 12 GeV proton synchrotron (PS) at High Energy Accelerator Research Organization (KEK) in Tsukuba, Japan. Protons extracted from PS hit an aluminum target. Secondary particles, mainly pions, are focused into Super-K direction by a set of magnetic horn system and decay into neutrinos in 200 m long decay volume. Properties of neutrino beam are measured at two points 250 km apart, with near detector (ND) system at KEK and with Super-K, a 50 kton water Čerenkov detector. The timing of accelerator extraction, 1.1 \( \mu \)sec every 2.2 sec, was synchronized to event timing at the far site using Global Positioning System (GPS). By using near detector data to characterize neutrino beam properties just after production, systematics related to e.g. uncertainties from absolute neutrino flux and absolute cross-section are greatly reduced.

Several instruments are placed in the beamline in order to monitor the beam quality\(^5\). Intensity and position of proton beam are measured by monitors located along the primary beamline. The kinematics of secondary pions just before the decay volume are measured using a specially-designed gas Čerenkov counter (pion monitor) in the early stage of the experiment. The intensity, profile and direction of tertiary beam are monitored by muon monitor system located just after the beam dump. Throughout the experiment, beam properties are stable and well within requirements not to affect oscillation measurements. For example, the stability of
beam direction has been better than 1 mrad, while requirement is 3 mrad. In total, $1.05 \times 10^{20}$ protons on target (POT) are delivered and data corresponding to $0.92 \times 10^{20}$ POT are used for the analysis.

2.2 $\nu_\mu$ disappearance final result

Events recorded with Super-K are selected with criteria similar to those with Super-K atmospheric neutrino analysis. We select events where the event vertex is inside fiducial volume and all the final state particles are fully contained inside the detector. Figure 2 shows the timing distribution of events observed in Super-K, after subtracting expected time of flight from KEK to Kamioka ($\Delta t$). In the left figure, open, hatched and shaded histograms are after removing pre-activity, requiring greater than $\sim$20 MeV energy deposit, and final selection, respectively. A clear peak is seen at $\Delta t = 0$, with negligible and expectation-consistent background from atmospheric neutrinos. Right figure shows zoom-up view around the expected timing (without fiducial volume requirement). Nine-bunch structure of KEK-PS is clearly seen, which gives us confidence that those events are due to neutrinos produced 250 km away.

![Figure 2: Timing distribution of events observed in Super-K.](image_url)

In the signal timing region of $-0.2 < \Delta t < 1.3 \mu$sec, 112 events are observed. From near detector measurements extrapolated to the far site, the expected number of events in case of no oscillation is $158.1^{+9.2}_{-8.6}$ events. For the near to far extrapolation, HARP hadron production data is used to constrain uncertainties from secondary pion kinematics. Further, for 38 events in which single Čerenkov ring identified as muon-like (against electron-like) is observed, we reconstruct the neutrino energy assuming two-body charged current quasielastic (CCQE) interaction, $\nu_\mu + n \rightarrow \mu + p$. Figure 3 (left) shows the reconstructed energy distribution, together with expectation without oscillation (dashed) and best-fit result with neutrino oscillation hypothesis (solid). The observed spectrum is consistent with oscillation best-fit parameters (Kolmogorov-Smirnov test probability of 37%) while not compatible with null oscillation (0.07%). We find evidence of neutrino oscillation from both reduction in number of events and distortion of energy spectrum.

Using both number of events and reconstructed energy spectrum, we performed maximum likelihood analysis to find best fit parameters of mixing angle ($\sin^2 2\theta_{23}$) and mass-squared difference ($\Delta m^2_{23}$). Best fit parameters are found to be ($\Delta m^2_{23}, \sin^2 2\theta_{23}$) = ($2.8 \times 10^{-3}$ eV$^2$, 1.0). Figure 3 (right) shows allowed regions of oscillation parameters with 68, 90, and 95% confidence level. Result from Super-K atmospheric neutrino L/E analysis is also shown. The null oscillation case is excluded with 4.3$\sigma$ significance from difference in the log likelihood. K2K has confirmed the neutrino oscillation result reported by Super-K. Recently, MINOS collaboration also confirmed K2K and Super-K results.
Figure 3: Left: Distribution of reconstructed neutrino energy. Points, solid line and dashed line are data, the best fit spectrum and expectation without oscillation, respectively. Right: Allowed region of neutrino oscillation parameters. Result from Super-K atmospheric neutrino L/E analysis is overlaid.

2.3 Search for $\nu_e$ appearance

Appearance of $\nu_e$ from $\nu_\mu$ beam is searched for by utilizing the excellent $e/\mu$ separation capability of water Čerenkov detector. Single ring, showering (electron-like) type events in Super-K are selected as candidates of $\nu_e$ charged current events.

The dominant background to this search is expected to be $\nu_\mu$ neutral current interaction, where only $\pi^0$ is generated ($\nu_\mu + N \rightarrow \nu_\mu + \pi^0 + N'$) and one of $\gamma$ from $\pi^0$ decay is missed in the reconstruction. Early result using about half of K2K data indicated that we needed better $\pi^0$ rejection efficiency to improve the sensitivity. Thus, we have developed a new algorithm for $\pi^0$ identification. A second gamma-ray ring candidate is reconstructed by comparison of the observed charge and expected light patterns calculated under the assumption that two showering rings exist. Thus, it always ‘finds’ the second ring in an event. We use the invariant mass of reconstructed $\pi^0$ to discriminate $\nu_e$ signal from $\pi^0$ background events. The performance of algorithm is verified using Super-K atmospheric neutrino data. The efficiency for expected $\nu_e$ signal is 70%, while 70% of $\pi^0$ background is rejected.

Figure 4: Left: remaining candidate for $\nu_e$ appearance search at K2K. Two rings reconstructed by $\pi^0$ identification algorithm are also shown. Right: Upper bound on $\nu_\mu \rightarrow \nu_e$ oscillation parameters at 90 and 99% confidence levels.

After all selection, we have found one candidate event, while 1.7 events of background are expected from MC simulation. Among them, 1.3 events are from $\nu_\mu$ interaction, dominated by NC $\pi^0$ production, and rest is from intrinsic $\nu_e$ contamination in the beam. Figure 4 (left) shows an event display of the remaining candidate. Because the observation is consistent with expected background, we derive excluded region on $\sin^2 2\theta_{13}$, where $\theta_{13}$ is mixing angle between $\nu_\mu$ and $\nu_e$ measured in $\nu_\mu \rightarrow \nu_e$ transition and $\sim \frac{1}{2} \sin^2 2\theta_{13}$, as a function of $\Delta m^2$. Figure 4
(right) shows the excluded region. At $\Delta m^2 = 2.8 \times 10^{-3}$ eV$^2$, upper limit on $\sin^2 2\theta_{\mu e}$ is set at 0.13 with 90% confidence level$^{11}$.

2.4 Neutrino-nucleus interaction studies

In addition to the oscillation analyses, we study neutrino-nucleus interaction using high statistics near detector data. Because we always need to interpret data using neutrino interaction model, it is highly important to understand the neutrino interaction with nuclei to perform neutrino oscillation experiments. We have published three interaction studies:

1. As described above, NC $\pi^0$ production is a main background for $\nu_e$ appearance search. The cross-section ratio of this mode to the total charged current is found to be $\sigma$(NC$\pi^0$) / $\sigma$(CC) = 0.064 ± 0.001(stat.) ± 0.007(syst.), consistent with our MC simulation$^{12}$. 

2. CCQE is the main signal for oscillation analysis and also used for normalization of other modes. The axial vector mass $M_A$, parameter to characterize the axial form factor in this interaction, is measured to be $M_A = 1.20 \pm 0.12$ GeV$^{13}$. 

3. From early stage of experiment, K2K had observed significant deficit in the forward scattering events, which limits the prediction accuracy of the neutrino energy spectrum at the far detector. The cross section of CC coherent pion production is found to be much less than the predicted value, which explained the observed discrepancy. The cross-section ratio to the total charged current is found to be $\sigma$(CC coherent $\pi$)/$\sigma$(CC) < 0.060 at 90% confidence level$^{14}$. 

Several other analyses are still ongoing and will be published in future.

3 SciBooNE – a bridge between K2K and T2K

One of K2K near detectors, SciBar$^{15}$ has been disassembled and shipped to U.S. for a new experiment at Fermilab, E954 (named SciBooNE)$^{16}$. The main goal of SciBooNE is measurements of neutrino-nucleus scattering cross sections below 1 GeV. Booster Neutrino Beamline at Fermilab, where currently MiniBooNE$^{17}$ is running, can provide a neutrino beam with very similar energy spectrum to that anticipated in T2K. Using well understood SciBar detector and intense, well understood beam from Booster Neutrino Beamline, precise study of neutrino interaction will be performed. In SciBooNE, run with anti-neutrino beam mode is also planned, which will help designing future CP violation study in long baseline experiments.

The experiment was approved by Fermilab PAC in December 2005. In summer 2006, detectors were moved from KEK to Fermilab and construction of the detector hall was started. As of April 2007, the hall construction has been finished and detectors are commissioned with cosmic ray. The experiment is expected to start data taking before summer 2007.

4 T2K – first 'superbeam' experiment

Based on the successful experience with K2K, T2K (Tokai-to-Kamioka) experiment$^{18}$ was proposed to further improve our knowledge on neutrino oscillation. The experiment was approved in 2003 and started the beamline construction in 2004. The primary goals of T2K are search for $\nu_\mu \rightarrow \nu_e$ oscillation with more than an order of magnitude better sensitivity than K2K, and precise measurement of $\nu_\mu$ oscillation parameters. The commissioning of beamline is planned to start in April 2009.
4.1 J-PARC facility

Intense neutrino beam will be produced by proton accelerators at Japan Proton Accelerator Research Complex (J-PARC) in Tokai village, Ibaraki prefecture. J-PARC accelerator system is designed to deliver 0.75 MW of beam power with fast extraction from its main ring. It consists mainly three parts: a linac, a 3 GeV rapid cycle synchrotron (RCS) and a main ring. In January 2007, the J-PARC linac was successfully commissioned. Commissioning of RCS and the main ring are scheduled in Japanese fiscal year 2007 and 2008, respectively.

4.2 Off-axis beam

In order to achieve the physics sensitivity we aim at T2K, it is indispensable to reduce background such as neutral current $\pi^0$ production for $\nu_e$ appearance search. Because the neutrino interaction cross section increases at high energy while the oscillation is expected to take place in rather low energy region of $< 1$ GeV, it is desired to reduce unwanted high energy tail while keeping as high intensity as possible in the signal region.

![Neutrino energy spectrum](image.png)

Figure 5: Left: Neutrino energy as a function of parent pion energy for various off-axis angle. Right: Neutrino energy spectrum for T2K beamline at various off-axis angle.

One of the characteristic feature of the T2K experiment is its 'off-axis' beam configuration. Figure 5 (left) shows the energy of neutrino emitted to certain angles with respect to the parent pion flight direction as a function of pion energy. At off-axis angles, the neutrino energy is rather independent of parent pion energy but dependent on the off-axis angle from the pion direction. Thus, by intentionally displacing the average direction of the secondary pions from the far detector direction, one can realize intense, narrow band beam tuned to the expected oscillation maximum at the far detector.

T2K is the first experiment to adopt this method, originally invented by BNL-ES89 group. Figure 5 (right) shows the expected neutrino spectrum at the far detector. Currently the off-axis angle at startup is planned to be 2.5°, although the beamline is designed to accommodate off-axis angles of 2–2.5°.

4.3 Beamline construction status

The construction of J-PARC neutrino beamline started in 2004. The primary beamline to transport protons to target was connected to the J-PARC main ring in Nov. 2006. The design and fabrication of beamline components, such as superconducting magnets, beam monitor, and power supply are on schedule.

The excavation of the target station, where the target and magnetic horns will sit, is ongoing. The upstream half of decay volume was already constructed. The target is made of graphite to withstand collision with intense proton beam. Design and test of the target cooling system is in progress. A set of three magnetic horns will be used to focus secondary particles. A prototype
of the first horn has been tested with the design current of 320 kA (cf: K2K horns were operated with 250 kA, NuMI horns with ~200 kA) for 850,000 pulses. Currently the third horn prototype is being configured to be included in the test. The design of beam dump is fixed and the first module of cooling core is assembled in March 2007.

4.4 Near and far detectors

The near detectors, consisting of two detector systems, will be placed about 280 m from the proton target. The ‘on-axis’ detector will monitor the beam direction, while ‘off-axis’ detector system will measure the neutrino flux, energy spectrum, and interaction cross-sections. The off-axis detector consists of several sub-detectors housed in a magnet, which will be reused from UA1/NOMAD experiments at CERN. The design of detectors is fixed and prototyping/test/fabrication is ongoing. The construction of detector hall will start from summer 2007.

The far detector, SuperKamiokande, recovered the original photo coverage and has been taking data from summer 2006.

4.5 Sensitivity to $\nu_\mu$ disappearance ($\Delta m^2_{23}$, $\theta_{23}$)

With the design intensity of 0.75 MW, we expect about 2,200 (1,600) $\nu_\mu$ (charged current) interactions in the fiducial volume of Super-K per a year if we had no neutrino oscillation. Using large statistics data optimized to the oscillation maximum, T2K aims to measure $\Delta m^2_{23}$ and $\theta_{23}$ with precisions of $< 10^{-4}$ eV$^2$ and 0.01, respectively.

4.6 Sensitivity to $\nu_e$ appearance ($\theta_{13}$)

With $5 \times 10^{21}$ protons on target of data ($\sim$5 years with design intensity), 103 $\nu_e$ appearance signal is expected for oscillation parameters of ($\Delta m^2_{23} = 2.5 \times 10^{-3}$, $\sin^2 2\theta_{13} = 0.1$). The expected background is 23, of which 13 is from intrinsic $\nu_e$ contamination and 10 from $\nu_\mu$ interaction, dominated by mis-identified NC $\pi^0$ production. The 90% confidence level (CL) sensitivity, assuming 10% uncertainty in background estimation, is shown in Fig. 6. At $\Delta m^2_{13} = 2.5 \times 10^{-3}$, the 90% CL limit of $\sin^2 2\theta_{13} = 0.008$.

![Figure 6: Left: Expected 68, 90 and 95% confidence level allowed region for $\nu_\mu$ disappearance parameters with $5 \times 10^{21}$ POT. Right: 90% confidence level sensitivity for $\nu_e$ appearance signal at T2K.](image)
5 Summary

We have successfully completed the K2K long baseline neutrino oscillation experiment. K2K confirmed $\nu_\mu$ oscillation reported by Super-K and set a limit to $\nu_\mu \to \nu_\mu$ oscillation in appearance mode. As the first experiment of this kind, K2K established validity and usefulness of long baseline experiment.

Based on this successful experience, the next generation long baseline experiment in Japan, T2K, is in preparation. The main goals of T2K are observation of $\nu_e$ appearance and precise measurement of $\nu_\mu$ oscillation parameters. Accelerator and beamline are being constructed. The design of near detector system is fixed and engineering design, prototyping and production of detector component are ongoing. The far detector, Super-K, recovered its full photo-coverage and is running. We expect the first beam for T2K in April 2009.

Acknowledgments

The author would like to acknowledge all the members of K2K and T2K collaborations for results presented here. He would like to thank organizers of this excellent conference for the invitation and great hospitality during the conference.

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MiniBooNE

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MiniBooNE is a short baseline neutrino experiment designed to confirm or refute the LSND observed excess of electron anti neutrinos in a muon anti neutrino beam. The experimental setup, data samples, and oscillation fit method are discussed. Although the result was not public at the time of the talk, MiniBooNE has since published results, which are discussed briefly as well.

1 Purpose

The Liquid Scintillator Neutrino Detector, or LSND, observed an excess of $87.9 \pm 22.4 \pm 6.0$ candidate electron anti neutrino events in a muon anti neutrino beam, consistent with a neutrino oscillation probability of $0.264 \pm 0.067 \pm 0.045\%$. The three independent observed mass splittings (atmospheric, solar, and LSND) cannot be explained in a three neutrino standard oscillation framework, and would require new physics explanations. The current favored solution to LSND would include additional 'sterile' neutrinos involved in oscillations. Given that the solar and atmospheric oscillations have been confirmed by multiple experiments, MiniBooNE's goal, then, is to confirm or refute $\nu_\mu$ to $\nu_e$ oscillations at high $\Delta m^2$. MiniBooNE has the same ratio of neutrino path to energy, or same probing of $\Delta m^2$ as LSND, and complements LSND with a different event signature and different systematics than LSND.

2 Experiment

2.1 Overview

The Fermilab Booster produces protons with kinetic energy at 8.89 GeV/c, and these are directed into a beryllium target placed inside a magnetic focusing horn. The subsequent mesons, predominantly $\pi^+$, are focused by the horn and decay to produce neutrinos. Past the decay region and 450 m of dirt sits a $\sim$ 1kton, mineral oil Cherenkov detector. This 12 m diameter
sphere has 1280 photomultiplier tubes (PMTs) on the inner region, providing 10% PMT coverage. The outer 'veto' region has 240 PMTs placed back to back, and these are used to reject the $\sim 10$kHz cosmic muon background during the beam spill.

2.2 Event Reconstruction

Events within the 1.6$\mu$s beam spill, with high enough PMT hits, and no substantial veto activity are neutrino candidate events. PMT hits distinct in time form 'subevents' in each neutrino interaction; a muon decaying in the tank, for example, produces two subevents, one each for the muon and the decay electron.

PMT hit topology, charge, and timing determine event types in MiniBooNE. In a charged current quasi-elastic interaction (CCQE) the incoming neutrino converts to its corresponding lepton partner. The outgoing lepton's flavor implies the flavor of the neutrino, and the Cherenkov ring observed indicates which lepton interacted. Muon events have a sharper Cherenkov cone because they are minimum ionizing particles; electrons and photons have broader rings due to scattering and showering. Basic quantities, such as charge, and timing of the hits, are used in conjunction with reconstructed event properties, such as track length of the lepton, the angle of the lepton with respect to the beam direction, to identify an event.

MiniBooNE's mean neutrino energy is approximately 700 MeV; roughly 40% of all interactions are CCQE. While only about 10% are neutral current single pion production (NC $\pi^0$), the $\pi^0$ can decay to two photons which appear as electron-like rings. Depending upon the topology of these events, they can be mis-reconstructed as electron neutrino events.

About 25% of the light in MiniBooNE is delayed, isotropic scintillation light. The amount of scintillation light as compared to the prompt Cherenkov light can give additional information distinguishing electron events from background.

3 Appearance Analysis

3.1 Electron neutrino selection

MiniBooNE uses two independent particle identification (PID) algorithms to select electron neutrino events: a likelihood analysis, and a boosted decision tree analysis. These two analyses use different reconstruction algorithms, oscillation fit code and methodology, and are susceptible to different sources of systematic errors.

A simple likelihood based analysis forms three PID variables. First, it compares the hits in the tank to an electron hypothesis and a muon hypothesis to form a PID variable distinguishing electron events from muon ones. Second, it compares hits to an electron hypothesis as compared to a $\pi^0$, or, a single electron-like ring to two electron like rings. Finally, cuts are applied to both of these variables and the output pion mass from the assumed two ring hypothesis.

The second method, boosted decision trees (BDT), is similar to a neural net. A decision tree takes a sample and applies a cut to a variable with the most signal to background separation possible. Then, it takes the second best variable to cut on, and cuts on it, and so forth. At each cut point, the sample is either cut on again, should more information be extracted, or no more cuts are applied, and the sample is a 'leaf'. If a leaf is predominantly signal, it is a signal leaf; background events on a signal leaf are called mis-classified events. Boosting is an additional method to separate signal from background. Mis-classified events are weighted more, and the tree remade. Hundreds or thousands of trees are produced, and then summed. Events on a signal leaf count as '+1', on a background leaf '-1', and the total gives a PID variable.
Table 1: Breakdown of events passing electron neutrino selection cuts (likelihood method) with systematic error in the signal region (reconstructed neutrino energy between 475 and 1250 MeV).

<table>
<thead>
<tr>
<th>Process</th>
<th>Number of events</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\nu_\mu$ CCQE</td>
<td>$10 \pm 2$</td>
</tr>
<tr>
<td>$\nu_\mu e \rightarrow \nu_\mu e$</td>
<td>$7 \pm 2$</td>
</tr>
<tr>
<td>Miscellaneous $\nu_\mu$ Events</td>
<td>$13 \pm 5$</td>
</tr>
<tr>
<td>$\text{NC } \pi^0$</td>
<td>$62 \pm 10$</td>
</tr>
<tr>
<td>$\text{NC } \Delta \rightarrow N\gamma$</td>
<td>$20 \pm 4$</td>
</tr>
<tr>
<td>$\text{NC Coherent } &amp; \text{ Radiative } \gamma$</td>
<td>$&lt;1$</td>
</tr>
<tr>
<td>Out of tank events</td>
<td>$17 \pm 3$</td>
</tr>
<tr>
<td>$\nu_e$ from $\mu$ decay</td>
<td>$132 \pm 10$</td>
</tr>
<tr>
<td>$\nu_e$ from $K^+$ decay</td>
<td>$71 \pm 26$</td>
</tr>
<tr>
<td>$\nu_e$ from $K^{0}_{L}$ decay</td>
<td>$23 \pm 7$</td>
</tr>
<tr>
<td>$\nu_e$ from $\pi$ decay</td>
<td>$3 \pm 1$</td>
</tr>
<tr>
<td><strong>Total Background</strong></td>
<td>$358 \pm 35$</td>
</tr>
<tr>
<td><strong>0.26% $\nu_\mu \rightarrow \nu_e$</strong></td>
<td><strong>163 \pm 21</strong></td>
</tr>
</tbody>
</table>

3.2 *Electron neutrino sample*

The $\nu_e$ appearance analysis selection cuts reduce a sample of over 100,000 neutrinos events down to 358 events, as shown in Table 1. The primary backgrounds are NC $\pi^0$, and the intrinsic electron neutrinos in the beam, from $\mu^+$ decay and kaon decay.

The rate of NC $\pi^0$ induced background is constrained by the NC $\pi^0$ events with two well-reconstructed photon rings. The measured rate for the 'clean' $\pi^0$ sample is compared to the simulation, and a reweighting factor determined in bins of $\pi^0$ momentum. This factor is then used to correct the predicted mis-reconstructed $\pi^0$ events in the $\nu_e$ sample.

3.3 *Sources of uncertainty*

The systematic errors included in the table cover primarily: flux, cross section and detector modeling uncertainties. In each case, MiniBooNE’s data or external measurements constrain the error.

HARP measured protons producing $\pi^+$ off beryllium at exactly MiniBooNE's beam energy. The differential cross section data from HARP is fit to a parameterization function, which is then used in the MiniBooNE beam simulation. For kaon production, external measurements were made with beams of energy spanning 9.5 GeV/c to 24 GeV/C, these are scaled to 8.9 GeV/c using a Feynman scaling model and then fit as well. Errors cover both the spread of the data as well as parameterization uncertainties.

The differential cross section for quasi-elastic scattering is measured from CCQE $\nu_\mu$ data. A fit to the shape of the four-momentum transfer ($Q^2$) distribution fixes an effective axial mass and nuclear effects parameter which is then applied to the $\nu_e$ CCQE sample.

In order to model light propagation in oil properly, we use a variety of internal and external measurements. The model includes: scintillation light (yield, spectrum, decay times), fluorescence (rate, spectrum, decay times), scattering (Rayleigh, Raman), absorption, reflection (off the tank walls and PMT faces) and PMT effects (single photoelectron charge response, charge linearity). External measurements such as scintillation light from the oil in a proton beam, from cosmic ray muons, fluorescence spectroscopy, time resolved spectroscopy and attenuation measurements of the mineral oil are also included. Finally, samples in MiniBooNE such as the cosmic ray muons, their decay electrons, and in-situ laser flasks constrain the model.
4 Oscillation Fit

Just as there are two parallel PID selection methods, the fit for oscillation was performed in two different ways. The likelihood analysis uses a CCQE $\nu_\mu$ sample to constrain the predicted intrinsic $\nu_e$ spectrum from muon decay and the predicted $\nu_e$ spectrum from $\nu_\mu$ oscillations. The BDT method performs a $\chi^2$ minimization fit between data and simulation for a $\chi^2$ that includes both $\nu_\mu$ and $\nu_e$ events along with their correlations. The $\nu_\mu$ data sample is used to reduce the size of the flux and cross section uncertainties. Much like a 'near to far' ratio, this cancels systematics which are the same for the two samples, and also reduces the $\nu_e$ uncertainties with the high statistics $\nu_\mu$ sample. The largest intrinsic $\nu_e$ sample comes from $\mu^+$ decay. As the parent $\pi^+$ decays to both the $\mu^+$ and $\nu_\mu$, and MiniBooNE subtends a small angle of the neutrino beam, the $\mu^+$ spectrum is closely related to the observed $\nu_\mu$ spectrum, and the additional knowledge of the $\mu^+$ spectrum limits what the $\nu_e$ from $\mu^+$ can be.

5 Result

Two weeks after the presentation at Moriond, the collaboration agreed to 'open the box', and unblind the $\nu_e$ sample. Less than a month later, the oscillation result paper was posted to the preprint server and submitted for publication\textsuperscript{5}. Although the presentation of this talk lacked any reference to the result, it is summarized here for completeness.

MiniBooNE did not observe an excess of events consistent with a two-neutrino oscillation explanation of the LSND observation. The final sensitivity is shown in Fig. 1. Within the main energy fit region of reconstructed neutrino energy between 475 MeV and 1275 MeV, the $\nu_e$ sample was consistent with no oscillation (see Fig. 2). However, an excess at lower than 475 MeV has been observed, but is still under investigation, and is not consistent with a simple oscillation model.

Acknowledgments

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References

Figure 1: Top: The final sensitivity curves of MiniBooNE in $\sin^2(2\theta) - \Delta m^2$ space within a two neutrino oscillation model. Black shows the MiniBooNE 90\% C.L., dash shows the sensitivity for the likelihood analysis. Blue shows the BDT analysis 90\% C.L. Bottom: MiniBooNE's 90\% C.L. is shown in solid black, along with KARMEN2 (dot) and Bugey (dash).
Figure 2: Events passing $\nu_e$ selection criterion as a function of reconstructed neutrino energy for the likelihood analysis. Top: Data is shown in black, expected background in solid red. Dashed shows the background with the best fit oscillation hypothesis. Intrinsic electron neutrino induced events are shown in solid green, and solid blue shows events from muon neutrinos. Bottom: Data, background subtracted shown in black, best fit oscillation shown in dashed red, solid green shows $\sin^2(2\theta) = 0.004, \Delta m^2 = 1.0 \text{eV}^2$, and solid purple shows $\sin^2(2\theta) = 0.2, \Delta m^2 = 0.1 \text{eV}^2$. 
MINOS RESULTS, PROGRESS AND FUTURE PROSPECTS

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The MINOS long baseline experiment has been collecting neutrino beam data since March 2005 and has accumulated $3 \times 10^{20}$ protons-on-target (POT) to date. MINOS uses Fermilab’s NuMI neutrino beam which is measured by two steel-scintillator tracking calorimeters, one at Fermilab and the other 735 km downstream, in northern Minnesota. By observing the oscillatory structure in the neutrino energy spectrum, MINOS can precisely measure the neutrino oscillation parameters in the atmospheric sector. From analysis of the first year of data, corresponding to $1.27 \times 10^{20}$ POT, these parameters were determined to be $|\Delta m_{32}^2| = 2.74^{+0.44}_{-0.26} \times 10^{-3}$ eV$^2$/c$^4$ and $\sin^2(2\theta_{32}) > 0.87$ (68% C.L.). MINOS is able to measure the neutrino velocity by comparing the arrival times of the neutrino beam in its two detectors. Using a total of 473 Far Detector events, $(\nu - c)/c = (5.1 \pm 2.9) \times 10^{-5}$ (68% C.L.) was measured. In addition, we report recent progress in the analysis of neutral current events and give an outline of experimental goals for the future.

1 Introduction

It is now well established that neutrinos have non-zero masses and that neutrinos mix. Their weak interaction eigenstates (or “flavour” eigenstates) $\nu_\alpha$ are related to their mass eigenstates $\nu_i$ by a unitary transformation $U$:

$$|\nu_\alpha\rangle = \sum_i U_{\alpha i}^* |\nu_i\rangle$$

(1)

$U$ is called the PMNS$^{1,2}$ matrix. Neutrinos are created and detected by weak interaction processes but their propagation in free space is described by their mass eigenstates, causing relative phases to change. This leads to the phenomenon of neutrino oscillations.

MINOS is a long baseline neutrino oscillation search based at FNAL. The neutrino beam created in the NuMI beamline is sampled in two locations, $\sim 1$ km from the beam production

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target underground at Fermilab, and again 735 km downstream in the Soudan Underground Laboratory, in Minnesota.

Using this setup, MINOS measures the oscillation parameters $|\Delta m^2_{32}|$ and $\sin^2(2\theta_{23})$ to world leading precision. In addition, MINOS searches for sub-dominant $\nu_\mu \rightarrow \nu_e$ oscillations, oscillations into sterile neutrinos, and $\nu \rightarrow \bar{\nu}$ transitions. The MINOS Far Detector can also be used to detect neutrinos created in the atmosphere.

2 Experimental setup

2.1 The NuMI beamline

A schematic of the NuMI neutrino beamline is shown in Figure 1. NuMI uses protons with a momentum of 120 GeV extracted from Fermilab’s Main Injector accelerator. They impinge a 95.4 cm long segmented graphite target producing secondary particles, mainly $\pi$ and $K$ mesons. The positively charged mesons are focused using a system of two pulsed, parabolic magnetic focusing elements, called “horns”. The secondaries subsequently decay in a 675 m long, 2 m diameter evacuated decay volume producing neutrinos.

Hadrons reaching the end of the decay volume without decaying are stopped in an beam absorber following the decay pipe. The beam absorber consists of a water-cooled aluminium core surrounded by a layer of steel blocks and an outer layer of concrete. The remaining muons from the meson decays are stopped in the $\sim 300$ m of rock that separate the MINOS Near Detector hall from the beam absorber.

The target position relative to the first horn and the horn current are variable. For most of the collection of data used in the oscillations analysis presented here, the target was inserted 50.4 cm into the first horn to maximize neutrino production in the 1-3 GeV energy range. The data described here were recorded in this position, between May 2005 and February 2006, and correspond to a total of $1.27 \times 10^{20}$ POT. The charged current (CC) neutrino event yields at the ND are predicted to be $92.9\% \nu_\mu, 5.8\% \bar{\nu}_\mu, 1.2\% \nu_e$ and $0.1\% \bar{\nu}_e$.

2.2 The MINOS detectors

The MINOS detectors are designed to be as similar as possible while operating in very different conditions. They are steel-scintillator sampling calorimeters, magnetized to $\sim 1.3$ T allowing to measure particle momentum using track curvature as well as range. The scintillator planes are divided into 4.1 cm wide strips rotated by 90° on subsequent planes to enable 3-dimensional event reconstruction. The scintillator light is captured by embedded wavelength shifting (WLS)
fibres and transported to the edge of the detectors, where the optical signal is converted using Hamamatsu M64 (Near Detector) and M16 (Far Detector) photomultiplier tubes.

Below 10 GeV, the hadronic energy resolution was measured to be $56\%/\sqrt{E[\text{GeV}]} \pm 2\%$ and the EM resolution was measured to be $21.4\%/\sqrt{E[\text{GeV}]} \pm 4.1%/E[\text{GeV}]$. The muon energy resolution $\Delta E_\mu/E_\mu$ varies smoothly from 6% for $E_\mu$ above 1 GeV where most tracks are contained and measured by range, to 13% at high energies, where the curvature measurement is primarily used.

3 Monte Carlo tuning

MINOS took data with 6 beam configurations obtained by varying the target position and the current in the focusing horns. The Monte Carlo simulations of neutrino fluxes strongly depend on the underlying models of hadron production in the target, which are presently poorly constrained at MINOS energies. The nominal simulations, based on FLUKA05, yielded energy spectra which did not match the high-statistics data in the Near Detector. A better agreement was achieved by smoothly adjusting the $p_z$ and $p_T$ of hadrons (mostly pions) produced off the graphite target. The resulting spectra, which describe the data much more closely than the nominal simulations, are shown in Figure 2.

Figure 2: Data and Monte Carlo energy spectra and their ratios are shown for six different beam configurations. The blue line corresponds to the untuned, the red line to the tuned hadron production model. The tuning significantly improves data/MC agreement in all cases.
4 \( \nu_\mu \) disappearance analysis

4.1 Event selection and extrapolation

In the Far Detector (FD), the signal from eight scintillator strips is read out by the same PMT pixel. Therefore, the initial step in the reconstruction of the FD data is the removal of the eightfold hit-to-strip ambiguity using information from both strip ends. In the ND, timing and spatial information is first used to separate individual neutrino interactions from the same spill. Subsequent reconstruction is done in the same way in both detectors. Tracks are found and fitted, and showers are reconstructed to be combined to events. For \( \nu_\mu \) CC events, the total reconstructed event energy is obtained by summing the muon energy and the visible energy of the hadronic system. To prevent human biases when assessing the oscillation analysis results, a blinding mechanism was applied to the FD data set. This procedure hid a substantial fraction of the FD events with the precise fraction and energy spectrum of the hidden sample unknown. Events are pre-selected in both detectors, by requiring total reconstructed energy below 30 GeV and a negatively charged track. The track vertex must be within a fiducial volume such that cosmic rays are rejected and the hadronic energy of the event is contained within the volume of the detector. The pre-selected \( \nu_\mu \) event sample is predominantly CC with a 8.6% neutral current (NC) event background estimated from Monte Carlo (MC) simulations. The fiducial mass of the FD and ND is 72.9% and 4.5% of the total detector mass respectively.

A particle identification parameter (PID) incorporating probability density functions for the event length, the fraction of energy contained in the track and the average track pulse height per plane provides separation of \( \nu_\mu \) CC and NC events. The PID is shown in Figure 3 for ND and FD data overlaid with simulations of NC and CC events. Events with PID above -0.2 (FD) and -0.1 (ND) are selected as being predominantly CC in origin. These values were optimized for both detectors such that the resulting purity of each sample is about 98%. The efficiencies for selecting \( \nu_\mu \) CC events in the fiducial volume with energy below 30 GeV are 74% (FD) and 67% (ND).

![Figure 3: Data and tuned MC predictions for the PID variable in the ND (top) and FD (bottom). The arrows depict the positions of the selection cuts. The FD MC distribution for CC events uses the best fit parameters discussed in the text.](image)

The measurement of the energy spectrum at the ND is used to predict the unoscillated spectrum at the FD. Fits to the ND data yield tuning parameters for the predicted neutrino flux. These fits are based on parameterisations of the secondary pion production at the NuMI target as a function of \( x_F \) and \( p_T \) as described in Section 3. The FD prediction must also take into account the ND and FD spectral differences that are present, even in the absence of oscillations,
due to pion decay kinematics and beamline geometry. This is achieved using the Beam Matrix method. It utilizes the beam simulation to derive a transfer matrix that relates the neutrinos in the two detectors via their parent hadrons. The ND reconstructed event energy spectrum is translated into a flux by first correcting for the simulated ND acceptance and then dividing by the calculated cross-sections for each energy bin. This flux is multiplied by the transfer matrix to yield the predicted, unoscillated FD flux. After an inverse correction for cross-section and FD acceptance, the predicted FD visible energy spectrum is obtained. The oscillation hypotheses are then tested relative to this prediction. A distinct extrapolation method, referred to as ND Fit was also applied to the data, yielding similar results.

In total, 215 events are observed below 30 GeV compared to $336.0 \pm 18.3 \text{(stat.)} \pm 14.4 \text{(syst.)}$ events expected in the absence of oscillations. The systematic error is most relevantly due to NC contamination, ND to FD normalization and the hadronic shower energy scale. In the region below 10 GeV, 122 events are observed compared to the expectation of 238.7$\pm$15.4$\pm$10.7. The observed energy spectrum is shown along with the predicted spectra for both extrapolation methods in Figure 4.

![Figure 4: Comparison of the Far Detector spectrum with predictions for no oscillations for both analysis methods and for oscillations with the best fit parameters from the Beam Matrix extrapolation method (left canvas). The estimated NC background is also shown. The last energy bin contains events between 18-30 GeV. The right canvas shows the ratio of data and best fit over the unoscillated predictions.](image)

### 4.2 Oscillation analysis

Under the assumption that the observed deficit is due to $\nu_\mu \rightarrow \nu_\tau$ oscillations, a $\chi^2$ fit is performed to the parameters $|\Delta m^2_{32}|$ and $\sin^2(2\theta_{23})$ using the expression for the $\nu_\mu$ survival probability:

$$P(\nu_\mu \rightarrow \nu_\mu) = 1 - \sin^2(2\theta_{23}) \sin^2 \left( \frac{\Delta m^2_{32} L}{4E} \right)$$

where $L$ is the distance from the target, $E$ is the neutrino energy, and $|\Delta m^2_{32}|$ is the atmospheric mass splitting. The fit included the systematic uncertainties mentioned above as nuisance parameters as well as the small contribution from selected $\nu_\tau$ events produced in the oscillation process. The resulting 68% and 90% confidence intervals are shown in Figure 5 as determined from $\Delta \chi^2 = 2.3$ and 4.6, respectively. The best fit parameter values are:

$$|\Delta m^2_{32}| = (2.74^{+0.44}_{-0.26}) \times 10^{-3} \text{eV}^2/c^4$$

and

$$\sin^2(2\theta_{23}) > 0.87$$

at 68% C.L. with a fit probability of 8.9%. At 90% C.L. $(2.31 < |\Delta m^2_{32}| < 3.43) \times 10^{-3} \text{eV}^2/c^4$, and $\sin^2(2\theta_{23}) > 0.78$. The data and best fit MC are shown in Figure 4.
If the fit is not constrained to be within the physical region, the best fit is at $|\Delta m_{32}^2| = 2.72 \times 10^{-3} \text{eV}^2/c^4$ and $\sin^2(2\theta_{23}) = 1.01$, with a decrease in $\chi^2$ of 0.2.

It is expected that the systematic uncertainties will be reduced with additional data. More details of this analysis are available in [11]. An update on this result using $2.58 \times 10^{20}$ POT is expected during the Summer 2007.

5 Neutrino time-of-flight analysis

MINOS uses GPS synchronized clocks to timestamp neutrino interactions in both detectors. This enables the measurement of the neutrino time-of-flight over a distance of 734 km and thus the determination of the neutrino velocity. Similar terrestrial experiments performed in the past used much shorter baselines of $\sim$ 500 m and higher beam energies ($> 30 \text{GeV}$).

The time of each PMT hit is recorded by the detector’s clock to a precision of 18.8 ns in the Near Detector and 1.6 ns in the Far Detector. The time of the earliest hit of each event is taken to be the time of the neutrino interaction. The interaction times $t_1$ and $t_2$ in the two detectors are corrected for known offsets and delays, which were determined using test stand measurements. In addition, the time of the beam extraction signal is subtracted from both times. The beam extraction signal has a fixed relation to the arrival of neutrinos in the MINOS detectors. All times are therefore measured relative to this reference.

The NuMI beam pulse is not instantaneous, but has a duration of 9.7 $\mu$s with a five-batch or six-batch intensity profile depending on the accelerator running mode. The Near Detectors measures this time-intensity profile with neutrino interactions to high precision. The measured time-intensity profile forms a probability density function which is folded with a Gaussian distribution with a width of $\sigma = 150 \text{ns}$ to account for the uncorrelated jitter of the two GPS clocks. The resulting distribution $P(t)$ describes the predicted arrival time distribution at the Far Detector (shown as a solid line in Figure 6 for the five- and six-batch modes separately).

For the time-of-flight measurement, 473 neutrino-induced events in the Far Detector were used. The time of each event was compared to the predicted arrival time distribution. The
Figure 6: Time distribution of PD events relative to prediction after fitting the time-of-flight. The top plot shows events in 5-batch spills, the bottom 6-batch spills. The solid lines show the normalized prediction curves.

Time-of-flight $\tau$ was found by maximizing an unbinned log-likelihood function:

$$L = \sum_i \ln P(t^i_2 - \tau).$$

(5)

The distribution of measured event times together with the predicted distribution for the best fit $\tau$ is shown in Figure 6. The time-of-flight of neutrinos was measured to be

$$2449.223 \pm 0.032(\text{stat.}) \pm 0.064(\text{syst.}) \mu s \quad 68\% \ \text{C.L.} \quad (6)$$

Comparing to the MINOS baseline this translates to

$$\frac{(v - c)}{c} = 5.1 \pm 2.9 \ (\text{stat.} + \text{sys.}) \times 10^{-5} \quad 68\% \ \text{C.L.} \quad (7)$$

The systematic error is due to uncertainties on the timing delays and offsets.

6 Progress in Neutral Current analyses

Analyses of neutral current neutrino interactions in MINOS are currently in progress. Neutral current interactions are interesting for several reasons. They form an important background to the $\nu_\mu \rightarrow \nu_\tau$ oscillation analysis described in these proceedings and their cross-sections at the energies relevant to MINOS are not very well known. Furthermore, an observed deficit of neutral current events at the Far Detector could be evidence for light sterile neutrinos.

Neutral current interactions are selected using three event quantities: the event length, the number of tracks in the event and the track extension. In events where both tracks and showers were found, the track extension measures how much longer a reconstructed track is compared to the hadronic shower it is accompanied by. Distributions of the selection variables in Near Detector data and Monte Carlo are shown in Figure 7 (left canvas).

Using these variables, an energy spectrum of neutral-current-like events was produced. This is shown in the right canvas of Figure 7. The Monte Carlo is shown as a red line with a systematic error band; the background due to wrongly selected charged current events is shown as a blue hatched distribution. Within the estimated systematic errors, due to flux, cross-section and energy scale uncertainties, data and Monte Carlo agree well.

The extrapolation of these Near Detector results to the Far Detector and a fit for oscillations to sterile neutrinos are currently being worked on. Results from this analysis are expected later this year.
7 Future Prospects

In addition to the analyses reported here, MINOS is pursuing a $\nu_\mu \rightarrow \nu_e$ oscillation analysis in order to measure or constrain the as yet unknown mixing angle $\theta_{13}$. With its expected final statistics, MINOS will potentially be able to improve on the current best limit from the CHOOZ $^{12}$ experiment.

Other areas of interest include appearance and disappearance measurements of anti-neutrinos in the Far Detector as well as several non-oscillation analyses using the large number of neutrino interactions in the Near Detector.

Acknowledgments

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References

THE ANTARES NEUTRINO TELESCOPE

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On behalf of the ANTARES collaboration

The ANTARES collaboration is building a high energy neutrino telescope in the Mediterranean sea, 40 km off the French coast. The complete detector will be a 3-dimensional array of 12 lines equipped with 900 photomultipliers, installed at a depth of about 2500 m. On March 2006, the first line was successfully connected to the deep underwater junction box. Since then data arrive at the shore station continuously. In September 2006 the second line was connected. By March 2007 the detector has successfully taken data with 3 lines and first neutrinos have been observed.

1 Introduction

The ANTARES collaboration has been set up in 1996. Today it involves groups from France, Spain, Italy, The Netherlands, Russia, Germany and Romania. From 1996 to 1999 an extensive R&D program had been successfully performed to prove the feasibility of the detector concept. The environment parameters at various deep sea sites have been studied and a deployment site of the experiment has been chosen. It is 10 km south of the Hyeres archipelago at 42° 50' N, 6° 10' E. It combines the advantage of the necessary depth with the vicinity to the coast and infrastructure (harbors of Toulon and La Seyne).

2 Detector design

After the successful R&D phase the construction of the ANTARES detector has been decided in 2000 and is progressing. The detector consists of 12 lines and a junction box which distributes the power and clock synchronization signals to the lines and collects the data. The junction box is connected to the shore by a 42 km electro-optical cable. The lines have an equipped vertical length of 350 m starting 100 m above sea floor. Their horizontal distance is about 65 m and they
are arranged to form a regular octagon on the sea floor. Each line is connected to the junction box with the help of a submarine using wet-mateable connectors. It is composed of 25 storeys with a vertical distance of 14.5 m. The lines are kept straight by the floating force of a buoy at the top and an anchor at the bottom. They float in the sea current and the positions of the active detector elements are permanently monitored by an acoustic calibration system.

Each storey (see Fig. 1) contains three 45° downward looking 10° photomultipliers inside pressure resistant glass spheres - the optical modules (OM). The electronic cards are inside a titanium cylinder at the center. Some of the storeys contain supplementary calibration equipment like acoustic or optical beacons.

The signals of each photomultiplier are readout by two ASICs. For simple pulses charge and arrival time are digitized and stored for transfer to the shore station. For more complex pulses the pulse shape can be digitized with a sampling frequency up to 1 GHz. The time stamps are synchronized by a clock signal which is sent in regular intervals from the shore to all electronic cards. The overall time resolution of the signal pulses is limited by the transition time spread of the photomultipliers ($\sim$1.3 nsec). All data are sent in the shore station. With a noise light rate of 70 kHz on the one photon level this produces a data flow of 1 Gbit/sec to the shore. In the shore station a PC farm performs a data filtering to reduce the data rate by at least a factor of 100.

3 Construction status

From November 1999 to June 2000 a "demonstrator line" had been operated to prove the feasibility of the foreseen project. In October 2001 the final electro-optical cable was deployed over a length of 42 km from the foreseen ANTARES site to La Seyne where the power station
and the control room are located. In December 2002, the junction box was connected to the remote cable end and was deployed. Since its deployment it is permanently monitored and has successfully operated. From March to June 2003 two test lines had been operated to validate all components for mass production of the full detector.

In March 2005 the first permanent line has been installed. It contains 3 storeys with optical modules and equipment to monitor the deep sea environment. Results from its first year of data taking have been published. One year later the first complete detector line equipped with 75 optical modules has been deployed and connected to the junction box. Since then a new line is completed about every two months. About every 6 months the new lines are connected to the junction box during a submarine campaign. The most recent connection operations happened in September 2006 and January 2007. Therefore in March 2007 the detector was operating with a total of 5 complete lines - 42% of its final size.

4 Physics performance

Most studies so far concentrated on charged current interactions of $\nu_\mu$:

$$\nu_\mu (\bar{\nu}_\mu) + N \rightarrow \mu^- (\mu^+) + X$$

(1)

The direction of the muon is reconstructed using the fact that it emits Cherenkov light under a well defined angle and does not suffer from multiple scattering at high energies. In ANTARES several reconstruction algorithms for muons have been developed. They use the direct Cherenkov hits but take also into account secondary effects like diffusion, dispersion and electromagnetic showers which accompany high energetic muons. This leads to an angular resolution for the muons of better than $0.2^\circ$ above 1 TeV for the above mentioned 1.3 nsec single pulse resolution. To obtain the neutrino angular resolution one has to consider also the interaction kinematics and gets $0.7^\circ$ at 1 TeV which decreases to the detector-dominated $0.2^\circ$ at 100 TeV.

The estimation of the neutrino energy is based on the measurement of the light output of the muon track in the vicinity of the detector. In the TeV range the light output increases with energy due to radiative processes. However the measurement is compromised by the facts that these radiative processes are stochastic, the neutrino interaction point is invisible in most cases and only a short fraction of the muon track is seen in the detector. Procedures have been developed which estimate the muon energy within a factor 3 for energies below 100 TeV and within a factor 2 for higher energies.

The effective area is another important parameter which characterizes the performance of the detector. Fig. 2 gives the effective areas for a neutrino flux before penetration of the Earth for
various nadir angle bins. This has the advantage that such an effective area can be folded directly with neutrino flux predictions from astronomical sources to obtain the number of expected signal events and it can be easily compared to effective areas of gamma ray telescopes. There are three differences with respect to the usually shown effective areas for muons: the overall scale changes from km$^2$ to m$^2$ due to the smallness of the neutrino cross section; the energy dependence becomes much stronger due to the almost linear raise of the neutrino cross section; the opacity of the Earth limits the effective area to values below 20 m$^2$.

Using the above performance parameters one can estimate that ANTARES will detect about 3000 upward going muon tracks from atmospheric neutrinos per year. They provide a detectable neutrino signal in the ANTARES detector, even during its construction phase.

5 Data taking

From March 2006 until September 2006 data have been taken with a single detector line. During this period the basic concepts of the trigger mechanisms and reconstruction algorithms could be validated. Due to the particular geometry of such a single line detector, it is mainly sensitive to close to vertical downward going atmospheric muons. An example event is shown in Fig. 3. Here the muon is recorded over the full length of 350m of the line. The arrival time of each photon is shown as function of height. Deviations of the real detector geometry from an exactly vertical shape are ignored in the figure. Nevertheless the measured points are well aligned, as expected for an exactly vertically downward going track and the slope refers to the speed of light in vacuum (300 m per 1000 nsec in the units of the plot).

After a short period with 2 operational detector lines late in 2006 there are 5 active lines at the beginning of February 2007 making ANTARES a truly three dimensional array. This allowed to test for the first time the calibration method on their full scale.

5.1 Calibrations

The detector lines move in the sea current. Therefore the position and orientation of all elements must be monitored with a time interval of a few minutes. For this purpose each line contains several acoustic receivers. They communicate with transmitters at the sea bottom. From each set of acoustic data the distance between the two elements can be established with a precision
of a few centimeters. Several such measurements, distributed in space, allow to determine their relative positions by triangulation. This global system was now successfully operated with the 5 line detector. Fine tuning of its parameters is ongoing to reach the desired precision of about 10 cm. The acoustic system is complemented by measurements of the inclination and twist of each storey which determine their orientation in space.

To verify the timing calibration of the detector elements and to monitor the water properties optical beacons are distributed in the detector array. They emit short but powerful light flashes which can illuminate neighbouring lines. Large amounts of data have been taken and are being analyzed.

5.2 Neutrino signal

Apart from calibrations the 5 line detector allows as well a full three dimensional reconstruction of muon tracks and the distinction of upward and downward going tracks. For the latter it is important to reduce the fraction of downward going muon tracks or muon bundles which are misinterpreted as upward going because the flux of downward going muons is several orders of magnitude more important than the upward going atmospheric neutrino flux. Such a method had been developed within a PhD thesis based exclusively on the analysis of simulated data. The left plot of Fig. 4 shows a track fit likelihood parameter which can be used to distinguish the two components. A cut of $\Lambda > -5.3$ is suggested to reduce the fraction of misinterpreted atmospheric muons in the sample of upward going tracks to less than 10%. Applying the same method to data without any additional tuning one obtains the right plot on Fig. 4. The similarity between the two plots is striking despite of some important differences: the real data have been taken with a 5 lines detector whereas the Monte Carlo study was done for the full detector; no positioning calibrations had been applied for the real data sample.

Applying the above mentioned cut on $\Lambda$ to this data sample three candidate events remain. These have been cross checked with an independent reconstruction method and by using an event display. One of them is illustrated in Fig. 5. The reconstructed upward going track passes in the vicinity of three detector lines leaving a large amplitude signal at each of them. The color code reveals that these signals are time ordered from bottom to top leaving no doubt on the upward going character of this track.
Figure 5: Neutrino candidate event from February 2007, the detected photons are distinguished by their arrival time (color: red: early, blue: late) and amplitude (size), the upward going muon from the neutrino charged current interaction (cyan line) has passed closely at three of the five lines making its identification as upward going track unambiguous.

6 Outlook

Data taking as well as the construction of the detector continue smoothly. In the near future ANTARES will be able to present more quantitative results on atmospheric neutrinos and other physics analyses. By January 2008 the construction of the full 12 lines detector will be completed.

References

RELIC SUPERNOVA NEUTRINOS

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Relic supernova neutrino detection is discussed, with particular emphasis on the Super-Kamiokande experiment. Presently under intensive study and discussion, a novel proposed modification to existing and future water Cherenkov detectors is presented. Enabling such detectors to identify neutrons will significantly enhance their capabilities for relic supernova neutrino detection as well as for a wide variety of other physics topics.

1 Introduction

On February 23rd of this year we celebrated a big anniversary — it has been twenty years since the observation of neutrinos from supernova SN1987A. But nearby supernovas are fairly rare events, occurring approximately once every few decades within our own galaxy. Consequently, the handful of neutrinos seen in 1987 remain the only neutrinos ever observed which originated from somewhere more distant than our own Sun.

However, there are many galaxies, and so on average one supernova explodes somewhere in our universe every second. All of the neutrinos which have ever been emitted by every supernova since the onset of stellar formation suffuse the universe. As a result, we are continuously bombarded with diffuse supernova neutrino background (DSNB) radiation. In fact, about 100,000 of these so-called “relic” supernova neutrinos pass through each one of us every second.

These neutrinos, if observable, could provide a steady stream of information about not only stellar collapse and nucleosynthesis but also on the evolving size, speed, and nature of the universe itself. This goal, until recently a distant dream of both theorists and experimentalists, may now be achievable within the next few years. The work which could make this possible, along with its far-reaching implications for other types of neutrino physics, will be the focus of this paper.
2 Know Your Limits

Super-Kamiokande (Super-K, SK) is the world’s largest underground water Cherenkov detector. With a total mass of 50 kilotons and a fiducial mass of 22.5 kilotons, it is located some 250 km west of Tokyo in an old zinc mine deep in the Japanese Alps. The detector is described in detail elsewhere.

In 2003 the Super-Kamiokande Collaboration conducted a search for supernova relic neutrinos based on the first five years of data from Super-K. Unfortunately, this study was strongly background limited, especially by the many low energy events below 18 MeV which swamped any possible DSNB signal in that most likely energy range. Consequently, this search could see no statistically significant excess of events and therefore was only able to set upper limits on the DSNB flux. These limits, which at 18 MeV were about one event per 22.5 kton per year per MeV, are just a bit higher than modern theoretical predictions of the flux. However, as a background-limited search a convincing DSNB signal would require at least another twenty years of Super-K data. Not a very encouraging situation!

But what if the DSNB events could be uniquely identified in SK, freeing the analysis from background complications?

3 How Can One Identify the DSNB Events?

In what began in 2002 as a search for a new method of extracting the relic supernova neutrino signal without background issues, theorist John Beacom and I have been tossing around ideas regarding modifying the Super-K detector. We finally decided to tackle the DSNB problem once and for all, with the admittedly ambitious goal of extracting a clear, positive signal within the next few years. It has proven to be a very fruitful partnership: in 2004 we published a Physical Review Letters article outlining our proposal, upon which this paper is largely based.

3.1 The Goal

All of the events in the present DSNB analysis are singles in both time and space. This singles rate is actually quite low in Super-K, only three events per fiducial ton per year. Therefore, if it were possible to look for coincident signals, i.e., for a positron’s Cherenkov light followed shortly and in the same spot by the gamma cascade of a captured neutron, then these troublesome background singles could be eliminated.

“Wouldn’t it be great if we could tag every supernova relic neutrino,” we thought. Well, the reaction we are looking for is:

$$\overline{\nu}_e + p \rightarrow e^+ + n$$

So the real question is, how can we reliably identify the neutron?

3.2 The Challenge

Of course, it is well known that free neutrons in water get captured by free protons and emit 2.2 MeV gammas, far below Super-K’s normal trigger threshold. However, if we could manage to see these we’d be in business! Maybe we could just lower the Super-K threshold briefly after each regular trigger...

This would be possible, and no SK change except a new trigger board would be required. Such a board has actually been built and tested in Super-K, but its efficiency for neutron detection is very low. As DSNB interactions within Super-K are rare, above 10 MeV just a handful per year, what we really want is to get all of this signal.
Hence, we need something in the water which will compete with the hydrogen in capturing neutrons.

3.3 The 0.1% Solution

We finally turned to the best neutron capture nucleus known: gadolinium. It has a nice 8.0 MeV gamma cascade, easily visible in Super-K. Unlike metallic Gd, the compound gadolinium (tri)chloride, GdCl$_3$, is highly water soluble.

We found that in order to collect 50% of the neutrons on gadolinium and 50% on hydrogen you’d need to put just 9 tons of GdCl$_3$ in Super-K. That’s exactly two cubic meters. No problem!

Even better, to collect >90% of the neutrons on gadolinium you’d only need to put 100 tons of GdCl$_3$ in Super-K. That’s about twenty cubic meters, or a 0.1% concentration of Gd in the tank, and with it we can tag almost all of the relic events.

Models vary, but with this solute in the water Super-K should see about five DSNB events each year with virtually no background at all. Now imagine a future, megaton-scale water Cherenkov detector like Hyper-Kamiokande (Hyper-K) observing 100+ supernova relic neutrinos every year...

3.4 The Price of Gd in China

From a physics standpoint it certainly seems like GdCl$_3$ is a nice compound to use for tagging neutrons, but can we afford 100 tons or more of it? As it turns out, there has been a dramatic revolution in the price of gadolinium over the past two decades. The opening of new mineral fields in the Gobi desert and the scaling up of rare earth refining and purification technologies have caused the price to plummet three orders of magnitude in recent years.

If we had tried to use gadolinium in Super-Kamiokande from day one the raw materials alone would have added $400 million dollars (U.S.) to the cost of that $100 million project. Today, acquiring 100 tons of 99.99% pure GdCl$_3$ will cost us just under $530,000. The formerly high price of gadolinium could very well explain why no one has ever even proposed using gadolinium in very large detectors before.

4 What Else Can We Do With Gd?

If adding GdCl$_3$ only allowed us to clearly see the as-yet-unobserved relic supernova neutrinos it would be a significant scientific breakthrough. In just a few years, the yield of supernova neutrinos from SN1987A would be obtained and then surpassed. Furthermore, since the cost of GdCl$_3$ is low this approach, unlike all previous neutron detection technologies, is readily scalable at reasonable expense to even the largest proposed future projects. The addition of gadolinium should add less than 1% to the total capital cost of any such experiment.

As it turns out, mixing GdCl$_3$ into Super-K’s water will open up a wide variety of new physics opportunities in addition to making possible the world’s first observation of the DSNB.

4.1 Galactic Supernovas

Naturally, if we can do relics, we can do a great job with galactic supernovas, too. With 0.1% gadolinium in the Super-K tank,

- the copious inverse betas get individually tagged, allowing us to study their spectrum and subtract them away from
- the directional elastic scatters, which will double our pointing accuracy.
The $^{16}$O NC events no longer sit on a large background and are hence individually identified, and

- the $O(\nu_e, e^-)$F events’ backwards scatter can be clearly seen, providing a measure of burst temperature and oscillation angle.

In addition, based on event timing alone, Super–K with GdCl$_3$ will be able to immediately identify a neutrino burst as a genuine supernova. This is due to the fact that the average timing separation between subsequent neutrino interactions would be much longer than the timing separation between coincident events (except for a very close supernova, but in that case see the “SN Early Warning” section). Even a modest number of these coincident inverse beta events would be a clear signature of a burst and could not be faked by mine blasting, spallation, or dropped wrenches.

These same distinctive inverse beta signatures will allow SK to look for black hole formation (and other interesting stuff) out to extremely long times after the burst. Above 6 MeV, coincident inverse beta background events, primarily due to the many nuclear power reactors in Japan, will occur on the level of less than one a day. This is to be compared with about 150 singles events a day in our final low energy sample. Therefore, the presence of Gd in the SK water will mean that signals from a supernova will take much longer to drop below the background level, making late neutrino observations of the cooling SN remnant possible for the first time.

4.2 SN Early Warning

Inspired in part by our PRL article’s preprint, another group of authors has pointed out the possibility of being able to tell that a wave of SN neutrinos was about to pass through the Earth:

Let’s suppose that a relative large, rather close star, like Betelgeuse, is about to explode as a supernova. Carbon burning takes about 300 years, then neon and oxygen burning each power the star for half a year or so. Finally, silicon ignites, forming an inert iron core. After about two days of Si burning, the star explodes as a supernova.

But during silicon burning the star is hot enough ($T > 10^9$ K) that the pair annihilation process

$$e^+ + e^- \rightarrow \nu_x + \bar{\nu}_x$$

starts to produce large numbers of $\nu_e$'s with an average energy of 1.87 MeV. This is coincidentally just above the inverse beta threshold of 1.8 MeV.

Therefore, if Super–K has GdCl$_3$ in it when this happens, we would expect to see ~ 1000 inverse beta neutron capture singles (the positron is not above Cherenkov threshold) a day. This is seven times the current low energy singles rate in SK, and could not be missed. No other detector on Earth would know that the main burst was about to arrive — only SK with Gd could do this! Surely the astronomical and neutrino communities, not to mention our gravity-wave colleagues, would appreciate knowing that a nearby star was about to explode.

Now, granted, the supernova has to be pretty close. This trick will only work well out to about 2 kiloparsecs in Super–K or 10 kpc in Hyper–K. On the other hand, these are the most valuable bursts and we would have the most to lose if we missed one due to calibration or scheduled detector downtime. Such downtime could be postponed a few days in the event of a sudden rise in the neutron capture rate. So, think of this as a supernova insurance policy.
4.3 Reactor Antineutrinos

If we were to introduce a 0.1% solution of gadolinium into Super-Kamiokande, we could collect enough reactor antineutrino data to reproduce KamLAND's first published results in just three days of operation. Their entire planned six-year data-taking run could be reproduced by Super-K with GdCl₃ in seven weeks, while Hyper-K with GdCl₃ could collect six KamLAND-years of $\bar{\nu}_e$ data in just one day.

Although Super-K with GdCl₃ will not be able to extract spectral information over the entire energy range to which scintillator detectors are sensitive, it will have the unique ability to provide some $\bar{\nu}_e$ directional information via the emitted positrons' Cherenkov light. This should, especially given the extremely high statistics involved, allow significantly tighter constraints to be placed on the solar neutrino oscillation parameters than any other method which could conceivably become operational before the close of the present decade, and possibly far beyond. We would have these data in hand within months of the decision to introduce GdCl₃ into Super-Kamiokande.

5 Gadzooks!

Since John and I were focusing on the low energy side of things, we haven't even gotten into how this solute should also allow our high energy friends to differentiate between atmospheric (or long baseline) neutrinos and antineutrinos of all species, reduce backgrounds to proton decay searches, and so on.

We propose calling this new project "GADZOOKS!" In addition to being an expression of surprise, here's what it stands for:

Gadolinium Antineutrino Detector Zealously Outperforming Old Kamiokande, Super!

6 Gadolinium R&D

But we never wanted to merely propose a new technique — we wanted to make it work!

In 2003 and again in 2005 and 2007 I received Advanced Detector Research Program grants from the U.S. Department of Energy for the study of GdCl₃'s properties and possible effects on Super-Kamiokande. These grants cover three main topics:

1. Explore the chemistry, stability, and optical properties of GdCl₃ in detail.

2. Understand any changes needed in the SK water system in order to recirculate clean water but not remove the GdCl₃ solute.

3. Soak samples of all materials which comprise the Super-K detector in water containing GdCl₃ over a period of greater than one year and then look for any GdCl₃-induced damage.

A scaled-down version of the Super-K water filtration system was built at the University of California, Irvine. We are currently using this system to test out new water filtration technologies in order to maintain the desired GdCl₃ concentration in otherwise pure water. Gadolinium retention rates of over 99.9% per pass have been achieved. Meanwhile, at Lawrence Livermore National Laboratory materials aging studies are underway. After a GdCl₃ exposure equal to 30 years at the proposed concentration in Super-K we see no significant damage to the aged detector components. Preliminary measurements of the optical properties of the GdCl₃ were conducted in Japan during the spring of 2004, with very promising results.

After two years of these bench tests, I was allowed to use the K2K experiment's one kiloton [KT] water Cherenkov tank, a 2% working scale model of Super-Kamiokande at KEK, for large-scale Gd studies. This was possible only after K2K's long-baseline neutrino beam turned off.
for good in early 2005 and final post-calibration runs were completed. In November of 2005 I introduced 200 kg of GdCl₃ into the KT.

The good news is that adding the gadolinium chloride itself did not hurt the water transparency in the KT tank, and the water filtering system developed at UCI worked perfectly. The bad news is that the chlorine attached to the gadolinium to make it dissolve in water attacked some old rust in the KT tank, which is made of painted iron, and lifted it into solution. This then made the water transparency go down and the water change color. Finally, at the end of March, 2006, we removed the GdCl₃ and drained the KT so we could look inside and be sure of what was happening.

This inspection of the inside of the KT tank showed large areas (about 20% of the total inner surface area) which had not been properly painted back in 1998 — these were very rusty. It is not believed that the GdCl₃ itself caused the rust. This has been checked with tabletop tests involving clean and pre-rusted iron samples soaked in GdCl₃ solutions. As Super–K is made of stainless steel, not (badly) painted iron, we still expect this idea will work in Super–K, though more studies are clearly needed.

It was decided that the next step in the gadolinium R&D should be to build a custom-made tank out of stainless steel, and make it as similar to Super–K as possible. In April, 2006, Lawrence Livermore National Lab agreed to fund the construction and operation of a stainless steel Gd-testing tank in the US. That study is now reaching completion, with results expected in mid-2007.

We learned a number of important things in the kiloton detector:

1. GdCl₃ is easy to dissolve in water.

2. The GdCl₃ itself (i.e., in the absence of old rust) does not significantly affect the light collection.

3. Choice of detector materials is critical with GdCl₃.

4. The 20-inch Super–K PMT’s operate well in conductive water.

5. Our Gd filtration system works as designed at 3.6 tons/hr and can easily be scaled up to higher (Super–K level) flows.

All of these findings are of course applicable to putting GdCl₃ into Super–Karnickande someday. Since Super–K is made of good quality stainless steel, not iron, we don’t expect such rust trouble there. Even so, we should (and will) make things work with gadolinium in a stainless steel test tank first.

It now appears quite likely that the decision will be made to put gadolinium into Super–K sometime in the next two years. The University of Tokyo is beginning to assign some of their young people to focus on the project, and we now have a Gadolinium Committee within the Super–K Collaboration - this is all extremely encouraging!

References

CUORICINO AND CUORE: BOLOMETRIC EXPERIMENTS FOR DOUBLE BETA DECAY RESEARCH


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The latest CUORICINO limit for the effective Majorana neutrino mass $\langle m_{ee}\rangle$ is presented. In addition, the sensitivity reach for CUORE, the next generation experiment currently under construction, and the background reduction schemes developed and tested are presented.

1 Introduction

The evidence of neutrino mass from oscillation and reactor experiments is one of the most important discoveries, obtained during the last years in the field of particle physics. Oscillation experiments can only measure the differences in the square of the masses. Still unknown are the absolute neutrino mass scale, their hierarchy, and whether they are Dirac or Majorana particles.

Double Beta Decay (DBD) is a rare transition, in which an even nucleus $(A,Z)$ decays in its isobar $(A,Z+2)$ with the emission of two electrons and two antineutrinos. The search for its neutrinoless variant (DBD$0\nu$) is a powerful and sensitive way to investigate the neutrino nature and mass. The next generation experiments have the possibility to reach down to a few meV scale. The signature for DBD$0\nu$ is a sharp peak in the energy spectrum at the Q-value of the transition due to the energy released by the two electrons emitted in the decay with no energy carried away by neutrinos. The measurement of the rate of the DBD$0\nu$ decay is related to the effective Majorana neutrino mass by $\Gamma_{0\nu} = G_{0\nu}(Q, Z)|M_{\text{nucl}}|^2 < m_{ee} >^2$, where $G_{0\nu}(Q, Z)$ is the phase space factor, $|M_{\text{nucl}}|$ are the nuclear matrix elements, and $\langle m_{ee}\rangle = \Sigma_k |U_{ek}|^2 m_k e^{i\phi_k}$, where $U_{ek}$ are the mixing matrix elements for the mass eigenstates $m_k$ and $\phi_k$ are the CP Majorana phases. A search for DBD$0\nu$ in several nuclei is imperative due to the theoretical uncertainties in $|M_{\text{nucl}}|$.

One of the promising nuclei is $^{136}$Te. Its high natural isotopic abundance of 34% alleviates the requirement for isotopic enrichment, and its high Q-value of 2530.3 keV results in a large phase space for the decay and lies in a relatively background free region between the Compton edge and the full peak of the $^{208}$Tl 2615 keV line. The large phase space contributes to a decay rate that is four or five times higher for $^{136}$Te than for $^{76}$Ge.

Two main approaches are used to search for DBD$0\nu$: homogeneous and non-homogeneous. In the non-homogeneous approach, an external source of the chosen DBD candidate is placed in the form of thin foils inside the detector. In the homogeneous approach, the detector material is chosen to be a compound containing the DBD candidate, providing a high efficiency and in many cases, high resolution technique.

2 CUORICINO

CUORICINO has a total TeO$_2$ mass of 40.2 kg consisting of an array of 62 TeO$_2$ crystals, arranged in a tower of 13 planes. 11 of the planes are made of 4 crystals of 5 cm cube, while 2 of the remaining planes have 9 crystals of 3x3x6 cm$^3$. CUORICINO uses a bolometric technique to search for DBD$0\nu$ of $^{136}$Te, in which an energy deposition in the crystals induces temperature increase. A measurable temperature increase can be achieved in dielectric and diamagnetic materials such as TeO$_2$ when operated at very low temperatures. The operating
temperature of CUORICINO is 10 mK. The temperature change of each crystal is detected using neutron transmutation doped Ge thermistors, thermally coupled to each crystal with glue spots.

The detector has a series of shields designed to reduce the background radioactivity. The inner-most lead shield is made of roman lead with 210Pb content of less than 4 mBq/kg (90% confidence level), 1.5 cm around the side and 10 cm on the top and the bottom of the tower. The background from the cryostat contamination is reduced by the copper shields, and 20 cm of commercial lead shielding and 10 cm of borated PET shielding surround the cryostat to reduce environmental gammas and neutrons in the detector volume.

2.1 Results

CUORICINO is currently running in the underground laboratory at LNGS (Gran Sasso National Laboratory). The data used for this analysis are from two separate runs, taken from April 2003 to September 2003 and from May 2004 to January 2007. The total statistics consists of 11.83 kg.y of 130Te. The average FWHM resolution at the 208Tl gamma line is 8 keV for the 5x5x5 cm$^3$ crystals, and 11 keV for the 3x3x6 cm$^3$ crystals.

The detected events are run through pulse-shape analysis for noise rejection, then through software filtering to optimize energy resolution. Coincident events were used to reduce and study the background. The detection efficiency for fully contained DBD0ν events is 86% as evaluated by Monte Carlo simulations. Rejecting coincident events reduces the background by roughly 20%.

The total measured background in the DBD0ν region is 0.18 +/− 0.01 cnts/keV/kg/y.

In evaluating the limit for the 130Te DBD0ν half-life ($T_{1/2}^{DB}$) the background spectra collected in the two runs are kept separate, due to the different measured FWHM and background. A maximum likelihood procedure is applied in the 2475 - 2550 keV energy interval to evaluate the maximum number of DBD0ν events possible with a flat continuum plus the 2505 keV 60Co gamma line. The Q-value for the DBD0ν peak is set at 2530.30 keV and a peak shape obtained by summing up N gaussians is considered, in order to account for the different detector energy resolutions. The obtained limit is of $3 \times 10^{24}$y at 90% C.L. (Fig. 7), corresponding to a limit for $\langle m_{\nu_e} \rangle$ between 0.16 and 0.81 eV (with nuclear matrix elements from 8). Limit variations on the order of 10% were observed when the expected peak was shifted by ±3 keV around the Q-value and for different models for the background fitting function and for the peak shape of the 60Co.

DBD0ν evidence for 76Ge have been claimed with a best value for $\langle m_{\nu_e} \rangle$ of 0.44 eV, in the degenerate region of the neutrino mass spectrum. Using the nuclear matrix element calculated by Klapdor, CUORICINO would set a 90% C.L. limit of 0.5 eV. The evaluated sensitivity for CUORICINO in 3 years live time, taken as the half-life corresponding to the minimal number
of detectable events above background at 1σ, is of $7 \times 10^{24}$, corresponding to $\langle m_{ee} \rangle \leq (0.1 + 0.6) \text{ eV}$. With this sensitivity, it is possible for CUORICINO to confirm the claimed evidence however, a null result will not rule it out due to the uncertainty in the nuclear matrix element calculations.

Next generation experiments, with $\sim 1$ ton of detector mass, and with a background in the DBD0ν region down to 0.01-0.001 $c/\text{keV/kg}/y$, are necessary to increase the sensitivity enough to reach the inverted hierarchy region of the neutrino mass spectrum.

2.2 Background analysis

An accurate knowledge of the radioactive sources, responsible of the background measured in CUORICINO in the DBD0ν region, is fundamental, in order to be able to reduce it to the wanted level in CUORE. The study of the coincidence and anticoincidence spectra, together with a comparison with Monte Carlo simulations of different radioactive contaminations in the bulk and surfaces of the various detector components, allowed us to identify the main sources of background in the DBD0ν region. They are β and α decays from $^{238}\text{U}$, $^{232}\text{Th}$ and $^{210}\text{Pb}$ contaminations on the crystal surface (20% ± 5%), α decays from the same contaminants on the surface of the mounting structure (50% ± 10%) and $^{208}\text{Tl}$ multi-Compton events due to $^{232}\text{Th}$ contaminations of the cryostat shields, far away from the detectors (30% ± 5%) (Fig. ??).

3 CUORE

The next generation experiment CUORE will be made of 19 CUORICINO-like towers, 52 $5 \times 5 \times 5$ cm$^3$ crystals each, arranged in a tight cylindrical structure (Fig. ??), for a total TeO$_2$ mass of $\sim 741$ kg. CUORE is expected to start at the beginning of 2011 with a sensitivity goal for $\langle m_{ee} \rangle$ down to the inverted hierarchy region of the neutrino mass spectrum. CUORE is expected to probe the range between 29 and 150 $\times t^{1/4}$ meV for a background level in the DBD0ν region of 0.01 $c/\text{keV/kg}/y$ and between 16 and 85 $\times t^{1/4}$ meV for a background level of 0.001 $c/\text{keV/kg}/y$. The high granularity of the detector will provide an effective reduction of the background by employing anticoincidence cuts. The detector shieldings have been designed in order to completely eliminate the contribution from environmental gamma radioactivity and from cryostat contaminations.

3.1 RED for background reduction

The technical feasibility of CUORE has been proven by the good performances of CUORICINO. The true challenge, as in all the next generation DBD0ν experiments, will be the background
achievement. The background knowledge acquired with CUORICINO was a helpful starting point for the background reduction R&D towards CUORE, aimed to reduce the background in the DBD0ν region to a level between 0.01 and 0.001 c/keV/kg/yr.

The shielding is designed to reduce the contribution from 208Tl multi-Compton events, due to 232Th sources in the cryostat shields, to a negligible level. The dominant background appears to be the surface contamination both of the crystals and the copper parts facing them. A “Radioactivity study Array Detector” (RAD), was built at the end of summer 2004. The detector consists of 2 planes of 5×5×5 cm³ TeO₂ crystals, with a structure almost identical to that of CUORICINO. The tests were performed in a second cryostat, housed in the Hall C of LNGS, provided with 5 cm thick internal copper shields and a 10 cm thick external lead shield. Due to the limited space in the cryostat, there is less shielding against the multi-Compton events from 208Tl in the DBD0ν region than in CUORICINO. This results in higher background in the DBD0ν region in the RAD runs and direct comparison between RAD run and CUORICINO is only possible above 3 MeV where no gamma lines from the U or Th are present. We were able to isolate the alpha decays occurring in the crystals bulk (peaks at the transition energy) and surfaces (broader and asymmetric peaks at the transition energy and at the alpha energy), and on the surfaces facing the crystals, whose largest one is due to copper (flat continuum from alpha energy down to low energies).

A series of RAD runs were performed. In the first run, the crystals were etched with nitric acid, removing about 10 microns of the surfaces, then polished with SiO₂ powder. The copper mounting structure was also etched and successively treated through electroerosion, removing from 10 to 20 microns of the surfaces. All the operations were performed in a clean environment. In the second run, all copper parts facing the crystals were fully covered with polyethylene film. The result was quite successful from the point of view of crystal cleaning: the TeO₂ surface contamination in 238U and 232Th was drastically reduced (a factor ~5), as proven by the reduction of the correspondent alpha peaks (Fig. ??). The extremely low background reached allowed us for the first time to disentangle the bulk vs. surface contamination of the crystals. Once the broader peaks due to surface contamination had disappeared, the Gaussian and sharp peaks due to crystals bulk contamination became visible. From the observed peaks and assuming secular equilibrium, TeO₂ bulk impurities have been evaluated: 232Th and 238U are present at a level of ~×10⁻¹³ g/g, ²¹⁰Pb concentration is of ~10⁻⁵ g/g. The background due to surface contamination of the copper was reduced by a factor 1.8 in the 3-4 MeV region by improved surface treatment and covering its surfaces with polyethylene film (Fig. ??).
Figure 4: Comparison between CUORICINO (black) and RAD (red) alpha background.

With the surface treatment techniques, material reduction and shielding developed in the R&D program towards CUORE, the dominant sources of background are expected to be crystal surfaces \(7 \times 10^{-3} \text{ c/keV/kg/y}\) and copper surfaces \(2.5 \times 10^{-2} \text{ c/keV/kg/y}\). A small contribution is expected from crystal bulk impurities \(10^{-4} \text{ c/keV/kg/y}\).

Additional R&D is underway to further reduce the background. One novel approach is the use "surface sensitive bolometers" (SSB) where the main bolometer crystal is completely surrounded by and thermally coupled to 6 thin bolometers. Surface events either from the crystal or the surrounding, would deposit energy on both the main crystal and a surrounding bolometer, and these events can be excluded via anticoincidence analysis. Unfortunately several wires were lost during our test run however, we were able to reduce the background by a factor of 2 between 3-4 MeV with respect to the first RAD. Better rejection is expected once all bolometers can be used for anticoincidence. Further tests are planned for 2008. This technique has the possibility to serve as a valuable R&D tool that allows us to unambiguously identify the origins of the background observed in CUORICINO.

In the meantime a new RAD test is going on in order to exclude possible contributions from thermal relaxations of the teflon spacers, used to clamp the crystals.

4 Conclusions

CUORICINO is running and sets a limit on the half-life of \(^{130}\text{Te}\) of \(3 \times 10^{24} \text{y}\) at 90\% C.L. and \(\langle m_{\text{ee}} \rangle\) between 0.16 and 0.18 eV. It demonstrates the technical feasibility of CUORE, the next generation experiment on the 1-ton scale, and has given us important insight into the sources of background. An R&D program is underway to achieve a background of below 0.01 counts/keV/kg/y.

Acknowledgments

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Prospect Neutrino Oscillation Measurements with Reactor Neutrinos

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The potential of neutrino oscillations measurements with reactor neutrinos is reviewed. This is the synthesis of the talk delivered during the XLII\textsuperscript{nd} Rencontres de Moriond: Electroweak Interactions and Unified Theories. La Thuile (Italy), March 2007.

1 Neutrino Oscillations

All experimental evidence\textsuperscript{1,2} (latest results\textsuperscript{3,4}), ranging many orders of magnitude of parameter space, strongly favours neutrino oscillations to be the dominant mechanism causing neutrino flavour mutations during neutrino propagation. Neutrino oscillations are the macroscopic manifestation of mixing in the leptonic sector. Neutrinos, therefore, interact as weak-force flavour neutrinos ($\nu_e$, $\nu_\mu$, $\nu_\tau$), while they oscillate during their propagation as mass neutrinos ($\nu_1$, $\nu_2$, $\nu_3$; since two different $\Delta m^2$ have been measured). The leptonic mixing is embodied by the so-called PMNS matrix causing a non-diagonal free Hamiltonian for neutrinos. For neutrino oscillations to happen, the neutrino mass spectrum must be non-degenerate, since the $\Delta m^2$ leads the oscillation phase factor as a function of $L/E$, as shown in the two $\nu$ equation:

$$P(\nu_\alpha \rightarrow \nu_\beta) = \sin^2 (2\theta) \sin^2 \left( \frac{1.27\Delta m^2 (\text{eV})^2 L (\text{km})}{E (\text{MeV})} \right)$$

The PMNS matrix can be parametrised in terms of 3 mixing angles ($\theta_{12}, \theta_{13}, \theta_{23}$) and a complex CP violating phase ($\delta_{CP}$). Additional Majorana phases (if present) are not observable via neutrino oscillations. Today (see Figure 1), $\theta_{12}$ and $\theta_{23}$ (dominating the so-called solar and atmospheric oscillation, respectively) are known to be large, $\theta_{13}$ is known to be small (dominating the effective decoupling between solar and atmospheric oscillations), while $\delta_{CP}$ is unknown. The leptonic CP violation is an important prediction, which, if measured, might be related to mechanisms contributing to the observed matter/antimatter asymmetry in the Universe\textsuperscript{5}. In order to measure the leptonic CP violation using neutrino oscillations, $\theta_{13}$ needs not to be zero\textsuperscript{6}. With the current sensitivity ($\sim$10% precision), no evidence for additional mechanisms beyond neutrino oscillations has been yet found. This possibility will be further explored by the next generation of high precision experiments characterising the PMNS matrix to unprecedented precision.

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Figure 1: Knowledge on $\theta_{12}$ (left) and $\theta_{13}$ (right) from a global analysis. $\theta_{12}$ is best measured by solar experiments and KamLAND, while the limits on $\theta_{13}$ are dominated by the CHOOZ experiment; although KamLAND and solar experiments also contribute. KamLAND inputs dominates our knowledge on $\Delta m^{2}_{\text{solar}}$, while MINOS and SuperKamiokande do likewise for the $\Delta m^{2}_{\text{atm}}$.

2 Reactor Neutrinos Oscillation Measurements

Reactor neutrinos can be used to carry out precision measurements of $\theta_{12}$ and $\theta_{13}$ mixing angles. Figure 2 shows the disappearance probability of reactor neutrinos, as appearance is impossible since there is no enough energy to create $\mu$ and $\tau$. Further precision on the knowledge of $\theta_{12}$ (measured in vacuum) is desired for comparison with the measurements obtained from solar neutrinos (matter driven) as well as improving the precision of calculations in the three neutrino scenarios. Reactor neutrino also allow clean measurement (if within reach) of the mixing angle $\theta_{13}$ necessary to measure $\delta_C P$. Sensitivities up to $\sin^2(2\theta_{13}) \sim 0.01$ can be reached by the forthcoming generation of liquid anti-neutrino multi-detector experiments.

Figure 2: Disappearance probability of reactor neutrinos versus baseline. Atmospheric (orange shaded) and solar (yellow shaded) oscillations parameters can be measured at baselines of $\sim 1$ km and $\sim 50$ km, respectively.
2.1 Virtues of Reactor Neutrinos

The inverse-$\beta$ reaction ($\bar{\nu}_e + p \rightarrow e^+ + n$) is a critical element of the experimental approach, granting a robust sample of neutrinos with very few analysis cuts. A neutrino is observed by detecting the correlated prompt-$e^+$ and delayed-$n$ energy depositions by a coincidence energy trigger. The $\nu$ energy can be accurately obtained from the $e^+$ energy for each event. Due to kinematics, there is an energy threshold: $\sim 1.8$ MeV (rather low). This threshold causes the sample of neutrinos detected ($\sim 25\%$ of total) to be those associated to fast decay (higher energy $\nu$) chains for fission products.

![Diagram showing the neutrino spectrum](image.png)

Figure 3: The neutrino spectrum visible (a): the reactor neutrino flux (b) and inverse-$\beta$ cross-section (c).

Figure 3 shows the reactor neutrino spectrum, which is the convolution of the sum over all the $\beta$-decay spectra of all fission debris and the inverse-$\beta$ cross-section (largest within energy regime and known within $\sim 0.2\%$)\(^7\). The active volume of this type of detectors consists of liquid scintillator (hydrogen rich) loaded with 0.1% of Gd, since the Gd thermal-neutron capture cross-section is very high granting an increase in detection efficiency. The reduction of the time coincidence ($\sim 8\times$) grants further background suppression compared with the case of neutron absorption on hydrogen. In addition, the neutron-to-Gd capture provides an energy tag: a cascade of gamma-rays amounting to $\sim 8$ MeV upon capture - well away from radioactivity singles. Therefore, even more background reduction is possible by requiring an energy cut to obtain a neutrino sample through the n-to-Gd capture.

A few other advantages are: i) reactor neutrinos are for free, ii) $e^+$ energy can be well calibrated with abundant nuclear sources on MeV range and has a lower limit ($>1.022$ MeV) useful to control systematics, iii) The oscillation baseline is precisely known and adequate to probe atmospheric and solar (unique) oscillations on Earth, iv) There is negligible neutral current contamination. Therefore, reactors provide high statistics and high precision neutrinos in $E/L$ necessary for the characterisation of the disappearance (always for reactor neutrinos) pattern used to infer the above mentioned measurements.
2.2 Backgrounds

Backgrounds arise from processes mimicking a time coincidence with n to Gd capture like emission. While a radio-pure detector is critical, all dominating backgrounds are associated to cosmic muons, hence, the overburden of each detector relates to its background rate. Therefore, each detector is position optimised upon their corresponding signal (distance to cores) to background (radio-purity and overburden) rate. There are 3 types of backgrounds: i) "accidental" caused by the random coincidence between natural radioactivity ($e^\pm$-like) and the capture of a thermal neutron on Gd (n-like) - or any event whose deposition exceeds 6 MeV. ii) "correlated" caused by an incoming fast-neutron, which first recoils on a proton ($e^\pm$-like) and, then, gets captured into Gd, once thermalised. iii) unstable spallation isotopes (generated on carbon) cause millisecons lifetimes $\beta-n$ decays, which are impossible to veto. The accidental background spectrum lies over oscillation analysis region; i.e. where the oscillation deficit is expected. This type of background can, however, be precisely measured in situ, letting the other contributions dominate the overall background uncertainties. Reactor-off data provides a measurement of the effective integrated background spectrum per detector. This is, however, a very rare occurrence, particularly unlikely in multi-core sites. So, alternative approaches to characterise backgrounds are mandatory. Background normalisation is generally poorly known, hence dedicated campaigns for their better understanding (at each site) will be carried out.

3 Solar Sector Measurement: $\sin^2(\theta_{12})$

In order to improve our knowledge on $\theta_{12}$ (large mixing) (Figure 2), unprecedented high statistics is desired. This need translates into the large size of hypothetical detector (labelled "New Reactor") able to yield integrated exposures $\sim 60$ GWkTy. The optimal location of the oscillation deficit associated to the first minima (for the neutrino spectrum maximum at $\sim 4$MeV) is $\sim 63$ km for $\Delta m^2_{\text{solar}} = 7.9 \times 10^{-5}$ eV$^2 \pm 4\%$.

![Figure 4: Precision on $\theta_{12}$ versus baseline by different potential experiments.](image)

As expected, the better the precision on the input knowledge of $\Delta m^2_{\text{solar}}$, the less stringent the dependence is on the optimal baseline. Figure 4 shows the precision on a measurement of $\sin^2(\theta_{12})$ (true value: $\sin^2(\theta_{12}) = 0.27$) versus neutrino propagation baseline for different experimental setups with different uncertainties in $\Delta m^2_{\text{solar}}$. Left shows the capability of the "New Reactor" in combination to further $\sim 3$ years of KamLAND (measuring $\Delta m^2_{\text{solar}}$ to $\sim 7\%$ at $3\sigma$). Right shows the "New Reactor" in combination with Super-Kamiokande doped with Gd$^3$ for 5 years exposure capable to measure $\Delta m^2_{\text{solar}}$ to $\sim 2\%$ (at $3\sigma$) within 5 years. In such
configurations, up to 2% (6%) at 1σ (3σ) precision can be achieved on $\sin^2(\theta_{13})$, assuming systematic uncertainties are controlled to 2%. Including the $\sin^2(2\theta_{13})$ uncertainty in the analysis deteriorates the sensitivity marginally to 3% (9%), respectively.

4 Atmospheric Sector Measurement: $\sin^2(2\theta_{13})$

In order to improve our knowledge on $\theta_{13}$ (small mixing) (Figure 2), unprecedented high precision is required. The location of the oscillation deficit associated to the first minima led by a non-zero value $\theta_{13}$ (for the spectrum maximum) is between 1 km and 1.5 km. The current limiting flux uncertainties (rate and spectral shape) known today to $\sim$2%\textsuperscript{10} will become negligible by the adoption of near detectors\textsuperscript{11}, whereby experiments will rely on the relative comparison between detectors spectra. Unlike beams, the near to far detector scaling is simpler: $\mathcal{O}(1/L^2)$.

![Figure 5: Evolution of the limit of $\sin^2(2\theta_{13})$ for $\Delta m^2_{3\ell} = 2.5 \times 10^{-3}$ eV$^2$.](image)

The strategy to achieve competitive measurements can be understood to a good approximation by the analysis of the sensitivity on $\sin^2(2\theta_{13})$ versus statistics available at the far detector as shown in Figure 5\textsuperscript{12}. A detailed description of the current scenario can be found in\textsuperscript{13}. Note there are 4 domains to be highlighted: i) (dashed-curve) the statistical limit improves with $1/\sqrt{N}$, achieved once the flux uncertainty is eliminated by near detector assuming no systematics. ii) (thick-curve) the statistical trend is limited by the inter-detector-normalisation uncertainty, causing a sort of plateau in the sensitivity. This trend is caused by the saturation of rate sensitivity dominating the limit at low statistics ($< 10^3$GW t years). iii) (thicker-curve) the plateau behaviour turns into a second statistical regime, as the sensitivity becomes shape dominated, independently of the inter-detector normalisation uncertainty. The shape becomes most relevant ($> 10^5$GW t years) as the bin-to-bin statistical power is large enough to resolve distortions caused by a non-vanishing value of $\theta_{13}$. iv) (thinner-curve) once shape uncertainties (i.e.}
“bin-to-bin”) are introduced, the sensitivity deviates from second statistical limit. Shape uncertainties arise from measured spectral differences between the near and far detectors caused, for example, by different background contributions or detector systematics. In summary, there are two strategies to measure $\sin^2(2\theta_{13})$ based on the statistics expected: a rate (measure integrated deficit) or a shape (oscillation spectral distortion) dominated measurement. The knowledge of the inter-detector normalisation (critical for the former) becomes, in theory, irrelevant for the latter in the very high statistics limit. In practice, currently there is no experiment can solely rely on the shape analysis, therefore, inter-detector normalisation is expected to be the most limiting uncertainty to most proposed experiments.

4.1 Detector Design

To achieve the fore mentioned control of systematics and backgrounds a new generation of $\theta_{13}$-LAND (liquid anti-neutrino detector) geared to yield accurate inter-detector comparisons has been developed. The far detector sets the statistics power of experiment, while the near detectors need not to be large (large statistics available). Therefore, the approach followed is to have many rather small ($< 20$ tonnes) identical detectors such that the detector-to-detector systematics - expected limiting uncertainty- can be more controlled. Additionally, using $N_{FD}$ far detectors has the extra advantage that each one can be regarded as an independent experiment, therefore, the overall far detectors uncorrelated uncertainties may scale down with $\sqrt{N_{FD}}$. Much of this trend has been set by the leading experiment in the field: Double Chooz\textsuperscript{14}. The RENO\textsuperscript{15} and Daya Bay\textsuperscript{16} collaborations have subscribed to such a trend, while future experiments such as Angra\textsuperscript{17} may opt for a single large $(500k)$ far detector. In addition, Daya Bay has designed their site (tunnels and laboratories) such that detectors can be swapped hoping to better understand detector systematics.

Figure 6: The Standard Detector $\theta_{13}$-LAND: The Double Chooz detector.

Figure 6 shows the 4 volumes standard detector: target (acrylics Gd loaded liquid scintillator), $\gamma$-catcher (acrylics unloaded liquid scintillator), buffer (non-scintillating oil) and inner-veto (water Čerenkov or liquid scintillator). There are also an outer inert shield to reject rock radiation and an active tracking outer-veto for cosmogenic background studies. The highlights
of this type of detector are: \(i\) hardware fiducial volume definition within acryllics (no position cut), \(ii\) low energy threshold below the \(e^+\) spectrum (\(\sim 0.5\text{MeV}\)), \(iii\) low singles rate within fiducial volume (< 10Bq), \(iv\) uniform detector response (no position cut and precise energy trigger), \(v\) possible in-coming fast-neutron tagging. The same scintillator batches and full readout system (\(10^7\) or \(8^7\) PMTs and electronics) are used in all detectors, minimizing so any detector response differences. Most effects are expected to be calibrated, monitored and corrected by multi-source (LEDs, lasers and radioactive sources) 3D calibration schemes. The use of the same calibration sources in all detectors minimizes the impact of the calibration uncertainties into the inter-detector normalisation; i.e. the absolute normalisation can, in principle, be relaxed. The detectors and calibration systems have been designed with a certain degree of information redundancy envisaged for the characterisation of possible systematic uncertainties.

![Figure 7: Double Chooz Sensitivity Evolution.](image)

4.2 The leading \(\theta_{13}\) Experiment: Double Chooz

The Double Chooz collaboration involves physicists from France, Germany, Japan, Russia, Spain, UK and US. Double Chooz is expected to be the first reactor experiment to have a word on \(\sin^2(2\theta_{13})\), whose vast R&D effort is summarised in\(^{14}\). The limited size of the far laboratory (former site of CHOOZ) limits the dimensions of the detectors (8.2t each and shown on Figure 6), while the reuse of this site allows an aggressive time-scale and valuable knowledge about backgrounds from CHOOZ. Figure 7 shows the improvement on the limit on \(\sin^2(2\theta_{13})\) for the far and the near-far detectors running. The latest estimation\(^{13}\) suggests that DC could reach 0.054 by 2010 and 0.028 by 2013, respectively.

The expected reach of the other collaborations realising a \(\theta_{13}\) reactor experiments (source\(^{13}\)): RENO \(^{15}\), Daya Bay \(^{16}\) and Angra \(^{17}\) claiming to reach sensitivities on \(\sin^2(2\theta_{13})\) up to 0.021 (by 2013), 0.009 (by 2015) and 0.006 (future), respectively. The KASKA \(^{18}\) collaboration joined Double Chooz, although they intend a next generation reactor neutrino. Daya Bay and Angra reach the high statistics necessary to partly exploit the shape information. The Hano-hano
concept was presented in the conference\textsuperscript{19}. A potential extension to Double Choz is hinted in\textsuperscript{12}.

5 Reactor-Beam Neutrinos Complementarity

Knowledge on $\sin^2(2\theta_{13})$ explored with reactor neutrinos can be used by beam experiments ($\nu_\mu$ beams). Beam experiments are sensitive to unknowns beyond $\theta_{13}$ (unlike reactor experiments), such as $\delta_{CP}$ and, in some cases, the sign of $\Delta m^2_{\text{atm}}$. The uncertainties associated to all PMNS parameters cause further ambiguity on the observable measured by a beam experiment, if an observation was made. Therefore, strong complementarity exists between results obtained by both beam and reactor based experiments\textsuperscript{20}. For example, if a reactor neutrino observation was made, the value of $\theta_{13}$ could be "fed" into the beam neutrino analyses to enhance their sensitivity on $\delta_{CP}$ by reducing the degenerate solution space. Therefore, it is worthwhile to pursue both approaches with comparable sensitivity for global analyses to infer the most about the structure of the PMNS mixing matrix.

6 Conclusions

Unique neutrino oscillation capabilities can be reached using reactor neutrinos. Capital measurements such as $\sin^2(2\theta_{13})$ (or limit) and $\sin^2(\theta_{12})$ under vacuum oscillations on Earth are within reach. New generation of LAND detectors has been developed to yield unprecedented systematic control specially within the detector-to-detector comparison systematics. Much of the corresponding R&D is already enriching the field of experimental neutrino and low background detection. Therefore, reactor neutrinos are still on the forefront of neutrino research and will provide complementary information to neutrino beams towards the better understanding of the leptonic sector on our Universe.

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POTENTIALITIES OF ATMOSPHERIC NEUTRINOS

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In this talk we will discuss the physics reach of the atmospheric neutrino data collected by a future megaton-class neutrino detector. After a general discussion of the potentialities of atmospheric neutrinos on general basis, we will consider concrete experimental setups and show that synergic effects exist between atmospheric and long-baseline neutrino data. Finally, we will show that present Super-Kamiokande data already have the capability to allow for a direct and unbiased measurement of the energy spectrum of the atmospheric neutrino fluxes.

1 Introduction

Despite their pioneering contribution to the discovery of neutrino oscillations, it is in general assumed that atmospheric neutrino data will no longer play an active role in neutrino physics in the coming years. This is mainly due to the large theoretical uncertainties arising from the poor knowledge of the atmospheric neutrino fluxes, which strongly contrast with the requirement of “precision” needed to further enhance our knowledge of the neutrino mass matrix. In this talk, we will show that despite these large uncertainties the atmospheric neutrino data collected by a megaton-class detector will still provide very useful information.

Let us begin by reviewing what we have learned so far about neutrino oscillations. The solar and atmospheric mass-squared differences are clearly determined, and we know that the solar angle is large but non-maximal while that the atmospheric angle is practically maximal. The present best-fit point and 1σ (3σ) ranges are:

\[
\begin{align*}
\theta_{12} &= 33.7 \pm 1.3 \left( ^{+4.3}_{-3.5} \right) , \\
\Delta m_{21}^2 &= 7.9^{+0.27}_{-0.28} \left( ^{+1.1}_{-0.89} \right) \times 10^{-5} \text{ eV}^2 , \\
\theta_{23} &= 43.3^{+4.3}_{-3.8} \left( ^{+0.8}_{-0.8} \right) , \\
|\Delta m_{31}^2| &= 2.6 \pm 0.2 \left( 0.6 \right) \times 10^{-3} \text{ eV}^2 , \\
\theta_{13} &= 0^{+5.2}_{-0.0} \left( ^{+11.5}_{-0.0} \right) , \\
\delta_{CP} &\in [0, 360] ,
\end{align*}
\]

leading to the following values for the elements of the leptonic mixing matrix, \( U \), at 90% CL:

\[
|U|_{\text{90\%}} = \begin{pmatrix}
0.81 & \cdots & 0.53 & \cdots & 0.00 & \cdots & 0.12 \\
0.32 & \cdots & 0.52 & \cdots & 0.60 & \cdots & 0.76 \\
0.27 & \cdots & 0.47 & \cdots & 0.65 & \cdots & 0.80 \\
\end{pmatrix} ,
\]

and at the 3σ level:

\[
|U|_{\text{3\sigma}} = \begin{pmatrix}
0.79 & \cdots & 0.50 & \cdots & 0.00 & \cdots & 0.20 \\
0.25 & \cdots & 0.47 & \cdots & 0.56 & \cdots & 0.79 \\
0.21 & \cdots & 0.42 & \cdots & 0.61 & \cdots & 0.83 \\
\end{pmatrix} .
\]
To set the basis for the following discussion, it is useful to study how much the information coming from atmospheric neutrino data contributes to this picture. The “solar” parameters $\theta_{12}$ and $\Delta m_{31}^2$ are completely determined by the solar and KamLAND data alone, and will therefore not be discussed here. As for the remaining parameters, in Fig. 1 we show the allowed regions implied by different combinations of neutrino experiments. It is clear from this figure that the bound on $\theta_{13}$ comes mainly from Chooz, with a small contribution from solar and KamLAND data, while the atmospheric mass-squared difference $\Delta m_{31}^2$ is mainly determined by the accelerator experiments K2K and Minos. The only parameter whose determination is still dominated by atmospheric data is $\theta_{23}$, being mainly a matter of total statistics. So it seems that indeed atmospheric experiments already have not much to say.

However, at a better look we note that present reactor and accelerator data (gray regions) exhibit a very high degree of symmetry. In particular, they have practically no dependence on $\delta_{CP}$, and they are totally insensitive to the neutrino mass hierarchy (sign of $\Delta m_{31}^2$) and to the octant (sign of $\theta_{23} - 45^\circ$). Conversely, when atmospheric data are also included in the fit these ambiguities, although far from being resolved, become at least non-symmetric. In particular, the atmospheric bound on $\theta_{13}$ (blue lines) considerably depends on the mass hierarchy, and the global fit with atmospheric data included (red lines) exhibit a weak but visible dependence on the CP phase and on the octant. Of course, this is not a proof that atmospheric data will be relevant in the future, especially since the results of forthcoming accelerator experiments such as T2K will have a non-trivial dependence on the value of $\delta_{CP}$ and on the neutrino mass hierarchy. However, it is a hint the very broad-range information provided by atmospheric neutrino data, which span about three orders of magnitude in length and more than five in energy, may still be complementary to accelerator experiments, which despite their high degree of precision are limited to a fixed baseline and cover only a very limited range in neutrino energy. In the rest of this talk we will present a systematic study of these potentialities.

2 Sensitivity to oscillation parameters

As already mentioned, the strength of atmospheric data is its very broad interval in neutrino energy ($E_\nu$) and baseline (determined by the nadir angle $\Theta_\nu$). In order to provide a global view of this whole range, in this section we will make extensive use of neutrino oscillograms of
the Earth, i.e. contours of equal probabilities in the neutrino energy-nadir angle plane. More exactly, we will show contours of equal oscillated-to-unoscillated ratio of atmospheric neutrino fluxes of flavor $\beta$ arriving at the detector, $\Phi_\beta / \Phi_\beta^0$, obtained by folding the primary neutrino flux $\Phi_\beta^0$ with the relevant conversion probability $P_{\alpha\beta}$, so that $\Phi_\beta = \sum_\alpha \Phi_\alpha^0 P_{\alpha\beta}$. As we will see, different regions of the $(E_\nu, \Theta_\nu)$ plane will show characteristic structures whose position and size is determined by various neutrino parameters.

$\theta_{13}$. Let us start by considering the sensitivity to $\theta_{13}$. As can be seen from Fig. 2(a), for non-zero value of $\theta_{13}$ matter effects induce a resonance in the $\nu_\mu \rightarrow \nu_e$ conversion probability at $E_\nu \sim 3 \div 6$ GeV. The precise position of this peak in the $(E_\nu, \Theta_\nu)$ plane depends on the value of $\theta_{13}$, so that in principle atmospheric neutrinos could be used to measure this angle as long as it is larger than about $\frac{\pi}{6}$. However, in practice the sensitivity is limited by two factors:

- **statistics**: at $E_\nu \sim 6$ GeV the atmospheric flux is already considerably suppressed;
- **background**: the $\nu_\mu \rightarrow \nu_e$ signal is diluted by the unavoidable background of $\nu_\mu \rightarrow \nu_e$ events.

Therefore, although some sensitivity is to be expected in a megaton detector, it is likely that atmospheric neutrinos will not be competitive with dedicated long-baseline and reactor experiments for what concerns the determination of $\theta_{13}$. However, an explicit observation of this resonance will provide a very important confirmation of the MSW and parametric-resonance mechanisms.

**Hierarchy.** As shown in Fig. 2(b), the sensitivity to the hierarchy arises from the observation of the same high-energy resonance which is also involved in the sensitivity to $\theta_{13}$. It is therefore only possible if $\theta_{13}$ is large enough. Note that in order to determine the hierarchy it is not sufficient to see the resonance (which would simply be a indication of non-zero $\theta_{13}$), it is also necessary to tell whether it occurred for neutrinos (normal hierarchy) or antineutrinos (inverted hierarchy). It is therefore particularly important to have a detector capable of charge discrimination. In the case of a non-magnetic detector such as a Water Cerenkov, if a resonance is observed it might still be possible to resolve the hierarchy by looking its size, since the number of neutrinos interacting in the detector is considerably larger than the number of antineutrinos and therefore
a normal (inverted) hierarchy would result in a larger (smaller) signal. Note, however, that the amplitude of the peak is affected both by $\theta_{13}$ and by the value of $\theta_{23}$, as we will see in the next paragraph, so that a poorly known value of these parameters will result in a considerable loss of sensitivity.

**Octant.** The sensitivity to the octant is one of the topics where atmospheric neutrino are mostly useful. As can be seen in Fig. 3(a), we have two characteristic signatures:

- **at low energy** ($E_\nu < 1$ GeV), we observe an excess (deficit) in the $\nu_\mu$ flux with respect to maximal mixing if $\theta_{23}$ is smaller (larger) than 45°. This effect is due to subleading oscillations induced by $\Delta m^2_{21}$, and is present also for $\theta_{13} = 0$. For $\theta_{13} \neq 0$ the $\nu_\mu$ flux arriving at the detector is modulated with the very fast $\Delta m^2_{31}$ oscillations, but the effect persists on average. Finally, this effect appears with the same sign for both neutrinos and antineutrinos, so that no charge discrimination is required for its identification.

- **at high energy** ($E_\nu > 3$ GeV), we observe a decrease (increase) in the $\nu_\mu$ flux with respect to maximal mixing if $\theta_{23}$ is smaller (larger) than 45°. This effect is again related to the matter resonance discussed for $\theta_{13}$ and the hierarchy, and indeed it appears only for $\theta_{13} \neq 0$. As already seen, in a detector without charge discrimination the signal could be considerably suppressed.

Note that the presence of a low-energy effect independent of $\theta_{13}$ guarantees a minimum sensitivity to the octant from atmospheric neutrinos, provided that the deviation of $\theta_{23}$ from maximal mixing is large enough. This is a unique feature which will prove very synergic with long-baseline data, as we will show in the next section. Note also that the slight preference for $\theta_{23} < 45°$ visible in Fig. 1 arises precisely from this effect, and from the observation of a small excess in sub-GeV $e$-like events in Super-Kamiokande data.

**CP phase.** Finally, let us spend a word on the sensitivity to the CP phase. A characteristic signal is visible in the intermediate-energy region, 1 GeV $< E_\nu < 3$ GeV, and arises from the interference of $\Delta m^2_{21}$-induced and $\Delta m^2_{31}$-induced oscillations. Although in principle it is observable, as Fig. 1 demonstrates with present data, this effect is quite small and probably
Figure 4: Allowed regions in $\sin^2 2\theta_{13}$ and $\delta_{CP}$ for LBL data alone (contour lines) and LBL+ATM data combined (colored regions). $^{3}$H$^{0}_{\text{ne}}$ and $^{3}$He$^{0}_{\text{ne}}$ refers to solutions with the true/wrong mass hierarchy and octant, respectively. The true parameter values are $\delta_{CP} = -0.85\pi$, $\sin^2 2\theta_{12} = 0.3$, $\sin^2 2\theta_{13} = 0.03$, $\sin^2 \theta_{23} = 0.6$, $\Delta m^2_{21} = 7.9 \times 10^{-5}$ eV$^2$, $\Delta m^2_{31} = 2.4 \times 10^{-3}$ eV$^2$.

It is hard to disentangle from other parameters. Moreover, its presence depends crucially on the size of $\theta_{13}$, so as for the sensitivity to the hierarchy it is not "guaranteed". In the fits which we will discuss in the rest of this talk the impact of $\delta_{CP}$ on the determination of the other parameters is properly taken into account, however we will not present any systematic study of the sensitivity of atmospheric data to $\delta_{CP}$ itself.

3 Synergies with long-baseline experiments

So far we have discussed the potentialities of atmospheric neutrinos in general terms. Let us now consider concrete experimental setups and compare their performances. In particular, we will focus on three proposed experiments:\(^3\)

- a Beta Beam ($\beta$B) from CERN to Fréjus (130 Km). We assume 5 years of $\nu_e$ from $^{18}\text{Ne}$ and 5 years of $\bar{\nu}_e$ from $^6\text{He}$ at $\gamma = 100$, with an average neutrino energy $\langle E_\nu \rangle = 400$ MeV. For the detector we assume the MEMPHYS Water-Cerenkov proposal, corresponding to 3 tanks of 145 Kton each;
- a Super Beam (SPL) from CERN to Fréjus (130 Km). We assume 2 years of $\nu_\mu$ and 8 years of $\bar{\nu}_\mu$ running, with an average energy $\langle E_\nu \rangle = 300$ MeV. Again we use MEMPHYS as detector;
- the T2K phase II (T2HK) experiment, corresponding to a 4MW super beam from Tokai to Kamioka (295 Km), with 2 years of $\nu_\mu$ and 8 years of $\bar{\nu}_\mu$. The detector is the proposed Hyper-Kamiokande, rescaled to 440 Kton for a fair comparison with the $\beta$B and the SPL.

A characteristic feature in the analysis of future LBL experiments is the presence of parameter degeneracies. Due to the inherent three-flavor structure of the oscillation probabilities, for a given experiment in general several disconnected regions in the multi-dimensional space of oscillation parameters will be present. Traditionally these degeneracies are referred to as follows:

- the intrinsic degeneracy: for a measurement based on the $\nu_\mu \to \nu_e$ oscillation probability for neutrinos and antineutrinos two disconnected solutions appear in the ($\delta_{CP}$, $\theta_{13}$) plane;
- the hierarchy degeneracy: the two solutions corresponding to the two signs of $\Delta m^2_{31}$ appear in general at different values of $\delta_{CP}$ and $\theta_{13}$;
- the octant degeneracy: since LBL experiments are sensitive mainly to $\sin^2 2\theta_{23}$ it is difficult to distinguish the two octants $\theta_{23} < 45^\circ$ and $\theta_{23} > 45^\circ$. Again, the solutions corresponding to $\theta_{23}$ and $\pi/2 - \theta_{23}$ appear in general at different values of $\delta_{CP}$ and $\theta_{13}$.
This leads to an eight-fold ambiguity in $\theta_{13}$ and $\delta_{CP}$, and hence degeneracies provide a serious limitation for the determination of $\theta_{13}$, $\delta_{CP}$ and the sign of $\Delta m_{31}^2$. In Fig. 4 we illustrate this problem for the $\beta$B, SPL and T2HK experiments. Assuming the true parameter values $\delta_{CP} = -0.85\pi$, $\sin^2 2\theta_{13} = 0.03$ and $\sin^2 2\theta_{23} = 0.6$ we show the allowed regions in the plane of $\sin^2 2\theta_{13}$ and $\delta_{CP}$ taking into account the solutions with the wrong hierarchy and the wrong octant. As visible in this figure, for the $\beta$B the intrinsic degeneracy cannot be resolved, due to the poor spectral information and the lack of precise information on $|\Delta m_{31}^2|$ and $\sin^2 2\theta_{23}$ (usually provided by the $\nu_\mu$ disappearance), while for the super beam experiments SPL and T2HK there is only a four-fold degeneracy related to the sign of $\Delta m_{31}^2$ and the octant of $\theta_{23}$. On the other hand, once atmospheric data are included in the fit all the degeneracies are nearly completely resolved, and the true solution is identified at 95% CL. This clearly show the presence of a synergy between atmospheric and long-baseline data: at least for this specific example, the combination of the two sets is much more powerful than the simple sum of each individual data sample.

To further investigate this synergy, in the left panels of Fig. 5 we show how the combination of ATM+LBL data leads to a non-trivial sensitivity to the neutrino mass hierarchy. For LBL data alone (dashed curves) there is practically no sensitivity for the CERN–MEMPHYS experiments (because of the very small matter effects due to the relatively short baseline), and the sensitivity of T2HK depends strongly on the true value of $\delta_{CP}$. However, by including the data from atmospheric neutrinos (solid curves) the mass hierarchy can be identified at 2$\sigma$ CL provided $\sin^2 2\theta_{13} \gtrsim 0.02 \div 0.03$. As an example we have chosen in that figure a true value of $\theta_{23} = 45^\circ$; in general the hierarchy sensitivity increases as $\theta_{23}$ increases. Note that the sensitivity to the neutrino mass hierarchy shown in Fig. 5 is significantly improved with respect to our previous results. There are two main reasons for this better performance: first, we use now much more bins in charged lepton energy for fully contained single-ring events; second, we implemented also the information from multi-ring events. This latter point is important since the relative contribution of neutrinos and antineutrinos is different for single-ring and multi-ring events. Therefore, combining both data sets allows to obtain a discrimination between neutrino and antineutrino events on a statistical basis. This in turn contains crucial information on the hierarchy, since as discussed in Sec. 2 the mass hierarchy determines whether the matter enhancement occurs for neutrinos or for antineutrinos.

In the right panel of Fig. 5 we show the potential of ATM+LBL data to exclude the octant degenerate solution. As seen in the previous section, this effect is based mainly on oscillations
with $\Delta m^2_{21}$ and therefore we have very good sensitivity even for $\theta_{13} = 0$; a finite value of $\theta_{13}$ in general improves the sensitivity. From the figure one can read off that atmospheric data alone can resolve the correct octant at 3$\sigma$ if $|\sin^2 \theta_{23} - 0.5| \gtrsim 0.085$. If atmospheric data is combined with the LBL data from SPL or T2HK there is sensitivity to the octant for $|\sin^2 \theta_{23} - 0.5| \gtrsim 0.05$. The improvement of the octant sensitivity with respect to previous analyses follows from changes in the analysis of sub-GeV atmospheric events, where now three bins in lepton momentum are used instead of one. Note that since in this figure we have assumed a true value of $\theta_{13} = 0$, combining the $\beta B$ with ATM does not improve the sensitivity with respect to atmospheric data alone.

4 Direct determination of atmospheric fluxes

So far we have discussed the potentialities of atmospheric neutrino for what concerns the determination of the neutrino parameters. The main message of the first part of this talk is that atmospheric data can provide very useful information on the neutrino mass matrix, despite the very large uncertainties in the neutrino fluxes. However, it is logically acceptable to invert the strategy, and to regard the poorly known atmospheric neutrino fluxes as a subject of investigation themselves. In this section we will therefore assume that the neutrino parameters have been accurately measured by other experiments, and we will show that it is possible to extract the atmospheric neutrino fluxes directly from the data.

There are several motivations for such direct determination of the atmospheric neutrino fluxes. First of all it would provide a cross-check of the standard flux calculations as well as of the size of the associated uncertainties (which, being mostly theoretical, are difficult to quantify). Second, a precise knowledge of the atmospheric neutrino flux is of importance for high energy neutrino telescopes, both because they are the main background and they are used for detector calibration. Finally, such program may quantitatively expand the physics potential of future atmospheric neutrino experiments. Technically, however, this program is challenged by the absence of a generic parametrization of the energy and angular functional dependence of the fluxes which is valid in all the range of energies where there is available data. We bypass this problem by using artificial neural networks as unbiased interpolants for the unknown neutrino fluxes. However, the precision of the available experimental data is not yet enough to allow for
a separate determination of the energy, zenith angle and flavor dependence of the atmospheric flux. Consequently in our work we have assumed the zenith and flavor dependence of the flux to be known with some precision and extract from the data only its energy dependence. Thus the neural network flux parametrization will be:

$$\Phi_{\alpha \pm}(E_\nu, c_\alpha, h) = F^{\text{fit}}(E_\nu) \Phi_{\alpha \pm}^{\text{ref}}(E_\nu, c_\alpha, h)$$  \hspace{1cm} (4)$$

where $F^{\text{fit}}(E_\nu)$ is the neural network output when the input is the neutrino energy $E_\nu$.

In Fig. 6 we show the results of our fit to the atmospheric neutrino flux as compared with the computations of the Honda and Bartol groups. The results of the neural network fit are shown in the form of a 1σ band, and plotted as a function of the neutrino energy. For comparison we also show the data from AMANDA. We see from this figure that the flux obtained from the fit is in reasonable agreement with the theoretical calculations, although the fit seems to prefer a slightly higher flux at higher energies. This indicates that until about $E_\nu \sim 1$ TeV we have a good understanding of the normalisation of the fluxes, and that the present accuracy from Super-Kamiokande neutrino data is comparable with the theoretical uncertainties from the numerical calculations. Note also that the results of our alternative fits depends only mildly on the choice of Honda or Bartol as the reference flux. This suggests that the present uncertainties on the angular dependence have been properly estimated, so that the assumed angular dependence has very little effect on the determination of the energy dependence of the fluxes. Thus the atmospheric neutrino flux determined with our method could be used as an alternative to the existing flux calculations.

5 Conclusions

In this talk we have discussed the potentialities of atmospheric neutrino data in the context of future neutrino experiments. We have shown that despite the large uncertainties in the neutrino fluxes atmospheric data will still provide useful information on the neutrino parameters, due to their very broad range in neutrino energy and nadir angle. In particular, we have proved that the sensitivity obtained by a combination of atmospheric and long-baseline data is much stronger than the one achievable by each data set separately. Finally, we have shown that present atmospheric data can be used to obtain a direct determination of the atmospheric neutrino fluxes.

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Sterile Neutrinos – neutrino masses, dark matter and baryon asymmetry –

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We discuss an extension of the Minimal Standard Model (MSM) by adding three right-handed neutrinos with Majorana masses being smaller than the weak scale or so, which we call as the \( \nu \)MSM. In this model, in addition to neutrino masses from oscillation experiments, dark matter and baryon asymmetry of the universe can be explained simultaneously. The lightest sterile (almost right-handed) neutrino with mass in keV range is a candidate of dark matter. The baryon asymmetry is originated in the oscillation between the heavier sterile neutrinos which are quasi-degenerate. In this talk, we review the basic features of the \( \nu \)MSM.

1 Introduction

The Minimal Standard Model (MSM) of particle physics has achieved a brilliant success in electroweak precision tests. In the last decade, however, there has been great progresses in experiments and observations, providing convincing evidence of physics beyond the MSM. Observations of non-vanishing neutrino masses in oscillation experiments conflict with the MSM, since neutrinos are exactly massless there. Moreover, observations of cosmic microwave background reveal that our universe is almost spatially flat and and is mainly composed of dark energy, dark matter and ordinary baryon. The MSM is incomplete to describe our universe.

The simplest way including neutrino masses is probably to add right-handed neutrinos \( \nu_R \). In this case, there appear neutrino masses, Dirac masses \( M_D \) as well as Majorana masses \( M_R \) for \( \nu_R \). Unfortunately, neutrinos masses in oscillation experiments are given in terms of \( M_D \) and \( M_R \) and their predictions are highly degenerate. Namely, the oscillation data do not tell us where are the scales of \( M_R \), which are mass parameters present in the Lagrangian.

A most popular one is the so-called seesaw scenario, where Majorana masses are superheavy \( M_R \sim 10^9 \text{–} 10^{15} \text{ GeV} \). In this case, the smallness of neutrino masses can be naturally explained by the seesaw mechanism \(^1\). Here neutrino Yukawa couplings are implicitly assumed to be comparable to those of quarks and other leptons. One may then expect that this fact and also
the origin of superheavy $M_R$ may be originated in grand unified theories. Furthermore, the baryon asymmetry of the universe can be generated via leptogenesis\textsuperscript{2}.

On the other hand, it has been proposed another possibility\textsuperscript{3,4}, where three right-handed neutrinos are introduced together with Majorana masses being smaller than the weak scale $\sim M_W$ or so. The model is called "the $\nu$MSM" as it is an extension of the MSM in neutrino sector. It has been shown that the $\nu$MSM can simultaneously explain not only neutrino masses and baryon asymmetry, but also dark matter of the universe\textsuperscript{3,4}. The $\nu$MSM introduces a new scale being much higher than $M_W$, and hence we may hope that the model will be verified experimentally. This point is a crucial difference between the $\nu$MSM and the seesaw scenario.

In this talk, we would like to explain some basic features of the $\nu$MSM. First, we discuss the dark matter in this framework and explain the current status of the dark matter scenario. Next, we see the baryogenesis scenario in the $\nu$MSM. The final section is devoted for summary.

2 What is the $\nu$MSM?

We start to explain the $\nu$MSM. It is an extension of the MSM by adding three right-handed neutrinos $\nu_{R,I}$ ($I = 1, 2, 3$) and its Lagrangian is given by

$$\mathcal{L}_{\nu\text{MSM}} = \mathcal{L}_{\text{MSM}} + i \bar{\nu}_{R,I} \gamma^\mu \partial_\mu \nu_{R,I} - F_{\alpha I} \bar{L}_\alpha \Phi \nu_{R,I} - \frac{M_I}{2} \bar{\nu}_{R,I} \nu_{R,I} + \text{h.c.},$$  \hspace{1cm} (1)

where $\Phi$ and $L_\alpha$ ($\alpha = e, \mu, \tau$) are the Higgs and lepton doublets, respectively. $F_{\alpha I}$ are neutrino Yukawa couplings and $M_I$ are Majorana masses. Here and hereafter we work in the basis where mass matrices of charged leptons and right-handed neutrinos are diagonal.

In this model, neutrinos obtain not only Dirac masses $M_D = F(\Phi)$ but also Majorana masses $M_I$. We assume the hierarchy between them, i.e., $[M_D]_{\alpha I} \ll M_I$. In this case, the seesaw mechanism can be applied, and neutrino mass eigenstates split into two groups. The lighter ones, called as active neutrinos $\nu_I$, obtain the mass matrix $M_\nu \simeq -M_D M_D^{-1} M_D^T$, which can be diagonalized by unitary matrix $U$. The others are called as sterile neutrinos $N_I$ with masses $M_I$. If we see the neutrino mixing in the charged currents, the left-handed neutrinos are written as $\nu_{L,I} = U_{\alpha I} \nu_I + \Theta_{\alpha I} N_I$, where $\Theta_{\alpha I} = [M_D]_{\alpha I} / M_I$ denotes the mixing matrix of active and sterile neutrinos and $|\Theta_{\alpha I}| \ll 1$ from our assumption. Thus, the neutrino oscillations can be explained by the flavor mixing among active neutrinos.

Importantly, we assume that Majorana masses are smaller than the weak scale $M_I < M_W$ in the $\nu$MSM. It should be noted that the seesaw mechanism still works even in our case. In the conventional seesaw scenario, the smallness of neutrino masses is a consequence of superheavy Majorana masses. In the $\nu$MSM it is due to the very small neutrino Yukawa couplings. As we will show, this smallness of the couplings is crucial to understand dark matter and baryon asymmetry in this framework.

3 Dark matter

A candidate of dark matter in the $\nu$MSM is the lightest sterile neutrino $N_1$ with mass $M_1 \sim$ keV, since it is charge neutral and is very long-lived. Note that active neutrinos cannot be dark matter. $N_1$ is not a completely stable particle, since there is no symmetry to forbid its decay. Indeed, for an interesting mass range, it mainly decays into three active neutrinos, and the lifetime is given by

$$\tau_{N_1} \simeq 5 \times 10^{26} \text{ sec} \left( \frac{\text{keV}}{M_1} \right)^5 \left( \frac{10^{-8}}{\theta^2} \right),$$  \hspace{1cm} (2)
Figure 1: The active-sterile mixing angle $\theta$ in terms of mass $M_1$. In the region between blue dotted and red dashed lines, the correct dark-matter abundance can be explained in the DW scenario. The regions excluded by various X-ray constraints are also shown.

where $\theta^2 = \sum_\alpha (|\Theta_{\alpha 1}|^2$. It can be seen that the lifetime of $N_1$ well exceeds the age of the universe ($\sim 10^{17}$ sec) in some parameter range.

Next, we consider the production of dark-matter sterile neutrino. The present abundance of $N_1$ should be consistent with the observational data, $\Omega_{\text{dm}}h^2 = 0.105^{+0.007}_{-0.013}$, where $\Omega_{\text{dm}}$ is the density parameter of dark matter and $h$ is the present Hubble constant in units of 100km/sec/Mpc. Due to the smallness of Yukawa couplings, $N_1$ is not thermalized in the early universe, and then it should be produced in non-thermal ways after the beginning of the hot big-bang universe (e.g., after the end of cosmic inflation).

The simplest scenario is that $N_1$ is produced solely from active-sterile neutrino transitions, which we call the Dodelson-Widrow (DW) scenario. In this case, the dominant production occurs at temperature $T_* \sim 100$ MeV ($M_1$/keV)$^{1/3}$. Recently, we have estimated the abundance of $N_1$ in the DW scenario. We have used the kinetic equation for density matrix of neutrinos, derived rigorously from the first principles of statistical mechanics and quantum field theory. Further, we have studied hadronic uncertainties in detail. Since $T_*$ is close to temperature of the quark-hadron transition, we cannot avoid uncertainties in the hadronic contributions to the production rate of $N_1$, and also to the QCD equation of state which is needed to determine the time-temperature relation. The detail discussions can be found in [7,8]. The result is shown in Fig. 1, where we show the parameter range of mixing angle $\theta$ and mass $M_1$ which give the correct dark-matter abundance in the DW scenario, including all theoretical uncertainties. We find that the required Dirac neutrino mass is typically $|M_D|_{\alpha 1} \sim 0.1$ eV.

We then consider cosmological constraints on the considering dark-matter scenario. One important constraint comes from X-ray observations. Note that $N_1$ can also decay into active neutrino and photon through one-loop diagrams with a small branching fraction. Emitted photons would be seen as a feature in X-ray background spectrum and also as a line X-ray spectrum coming from, clusters, galaxies, dwarf galaxies, etc. Since no signal has been observed, we obtain the upper bounds on the mixing angle $\theta$. The bounds are also shown in Fig. 1. It is seen that the X-ray constraint puts the upper bound on mass in the DW scenario. We find $M_1 < 6$ keV, as the most conservative bound, when all hadronic uncertainties are taking into account and the most stringent X-ray bounds are relaxed by a factor of two due to intrinsic uncertainties.
Interestingly, the lower bounds on mass are also obtained from cosmological constraints. The Tremaine-Gunn bound\textsuperscript{11} on fermionic dark matter applied to the dwarf spheroidal galaxies gives $M_1 > (\langle |q_a| \rangle / \langle |q_{ab}| \rangle) M_0$. From the matter power spectrum inferred from the Ly-$\alpha$ forest data, $M_0 \simeq 14.4$ keV\textsuperscript{13} and $M_0 \simeq 10$ keV\textsuperscript{14}. We find that $\langle |q_a| \rangle / \langle |q_{ab}| \rangle \simeq 0.8$ for $M_1 \sim 10$ keV\textsuperscript{8}, which leads to the lower bounds $M_1 \gtrsim 11.6$ keV and 8 keV, respectively. Therefore, the current cosmological constraints rule out the DW scenario in spite of all theoretical uncertainties involved in the estimation of the $N_1$ abundance. Note that, even if the decays of heavier sterile neutrinos in the $\nu$MSM cause the entropy production\textsuperscript{15}, this conclusion does not change\textsuperscript{8}.

We should stress that the sterile neutrino dark-matter is still possible if one goes beyond the simplest DW scenario. One example is the resonant production due to large lepton asymmetry of the universe\textsuperscript{16}. Another is to introduce new interactions of sterile neutrinos, e.g., $N_1$ may be produced in decays of some new scalar bosons\textsuperscript{17,18}. When such other production processes of $N_1$ are more effective, we may avoid the above cosmological constraints.

Finally, we comment on the implication to neutrino physics. We have found Dirac masses should be $|M_D|_{ab} \lesssim 0.1$ eV to avoid the overclosure by $N_1$ in any case. This means that the lightest sterile neutrino is irrelevant for masses and mixings in neutrino oscillations. We then need at least three right-handed neutrinos to explain oscillation data and dark matter at the same time\textsuperscript{3}. Further, when the number of $\nu_R$ is three as in the $\nu$MSM, mass of the lightest active neutrino is predicted as $m_1 \lesssim 6.7 \times 10^{-5}$ eV, which excludes a possibility that three active neutrinos are degenerate in mass\textsuperscript{3,8}.

4 Baryon asymmetry

We turn to discuss the baryon asymmetry of the universe in the $\nu$MSM. Let us see how the baryogenesis conditions\textsuperscript{19} can be satisfied in the $\nu$MSM. First, the baryon number $B$ and the lepton number $L$ are broken by the electroweak anomaly as in the MSM, and the sphaleron processes with $B + L$ violation (while $B - L$ conservation) are in thermal equilibrium for temperatures $T \sim 10^2 - 10^{12}$ GeV\textsuperscript{20}. Note that $L$ is also broken by non-zero Majorana masses in the $\nu$MSM, but its effect becomes negligible for high temperatures $T \gg M_W$ since $M_L < M_W$. Second, the $\nu$MSM contains 6 CP-violating phases in the lepton sector in addition to 1 CP phase in quark sector. Finally, the $\nu$MSM cannot realize the electroweak phase transition of strong first-order nature, taking into account the experimental bound on Higgs mass, as in the case of the MSM\textsuperscript{21}, and hence the electroweak baryogenesis cannot be realized. However, due to the smallness of neutrino Yukawa couplings, sterile neutrinos can be out of equilibrium in the $\nu$MSM. This point is crucially different from the case in the MSM. Therefore, the $\nu$MSM can satisfy all the baryogenesis conditions.

It has been proposed\textsuperscript{22} that the baryon asymmetry can be originated in sterile neutrino oscillations. The basic idea is as follows: Sterile neutrinos are created and oscillate together with CP violations. The total lepton number $\Delta L_{tot}$, defined by the sum of the lepton number of left-handed (active) lepton $\Delta L = \sum_\alpha \Delta L_\alpha$ and the helicity number of right-handed (sterile) neutrino $\Delta N = \sum_j \Delta N_j$, is not generated and is zero essentially due to the smallness of Majorana masses. However, the total lepton number can be distributed between active and sterile sectors. Namely, the asymmetries in the active leptons and sterile neutrinos can be generated such that their sum is still zero. Finally, the asymmetry stored in the active sector is partially converted into the baryon asymmetry due to the sphaleron processes.
Physics of this baryogenesis scenario is rather complicated and involved; the production and oscillation of sterile neutrinos, the generation of asymmetry in each lepton flavor and the transfer of asymmetries between active and sterile sectors. We have to take into account the coherent evolution of sterile neutrino numbers and also the decoherent effects due to the interaction with the background thermal plasma. We have investigated these issues in $^4$ and the kinetic equation for density matrix describing these phenomena have been presented there. Here we do not present it due to the limitation of space, but just explain the key steps in the generation of asymmetries. At the order $O(F^2)$, heavier two sterile neutrinos are produced in scatterings with top quarks and start to oscillate of each other. The CP violating effects in these processes are suppressed due to the constraints realizing the sterile neutrino dark-matter. At the order $O(F^4)$, the flavor asymmetries of active leptons, $\Delta L_{\alpha}$, are generated, but the total asymmetry of active leptons are conserving $\Delta L_{\alpha} + \Delta L_\mu + \Delta_\tau = 0$. These asymmetries arise from the fact that the propagation rates of active lepton and its anti-particle are different in the background with the sterile neutrino oscillation because of the CP violation in the lepton sector. At the order $O(F^6)$, the total asymmetry of active leptons is generated, i.e., $\Delta L = \sum_{\alpha} \Delta L_{\alpha} \neq 0$. This is because each flavor asymmetry $\Delta L_{\alpha}$ evolves differently due to the difference of Yukawa couplings. Notice that $\Delta L_{\alpha} = \Delta L + \Delta N = 0$ as mentioned before and then $\Delta N$ is also generated at this order. These asymmetries start to be generated at the temperature $T_L \sim (\Delta M_{32}^2 M_p)^{1/3}$ where $\Delta M_{32}^2 = M_3^2 - M_2^2$ and $M_p$ is the Planck scale, i.e., at the time when the effect of the sterile-neutrino oscillation becomes important.

The final result of the baryon asymmetry of the universe, the baryon to entropy ratio, is given by $^4$

$$\frac{n_B}{s} \approx 2 \times 10^{-10} \delta_{\text{CP}} \left( \frac{10^{-5}}{\Delta M_{23}^2 / M_2^2} \right)^{2/3} \left( \frac{M_2}{10\text{GeV}} \right)^{5/3},$$

where $\delta_{\text{CP}}$ is an effective CP violating parameter which is written in terms of mixing angles and CP phases in the neutrino Yukawa matrix $F$. $\delta_{\text{CP}}$ can be large as $O(1)$ being consistent with the oscillation data. The observational date $n_B/s \simeq (8.8-9.8) \times 10^{-11}$ can be explained by heavier sterile neutrinos which masses are quasi-degenerate. The mass degeneracy is bounded from below since $T_L \gtrsim M_W$. Otherwise, the total asymmetry of active leptons cannot be converted into baryon asymmetry via the sphaleron processes. The successful scenario restricts the mass range of quasi-degenerate sterile neutrinos as 17 GeV $\gtrsim M_{2,3} \gtrsim 1$ GeV being consistent with oscillation experiments. Here, heavier sterile neutrinos are in thermal equilibrium for $T \gtrsim M_W$ when $M_{2,3} \gtrsim 17$ GeV and the generation of baryon asymmetry is strongly suppressed. On the other hand, when $M_{2,3} \lesssim 1$ GeV, heavier sterile neutrinos decay after the big-bang nucleosynthesis and spoil its success. Quite interestingly, the above mass range is consistent with what we have assumed, all Majorana masses are smaller than $M_W$ or so. Therefore, the $\nu$MSM, being different from the MSM, can explain the origin of the baryon asymmetry of the universe.

5 Summary

We have discussed the extension of the MSM by adding three right-handed neutrinos with Majorana masses being smaller than the weak scale $\sim M_W$ or so, i.e., the $\nu$MSM. The model can solve the problems in the MSM, namely, the problems of neutrino masses, dark matter and baryon asymmetry of the universe at the same time.

In this model, the lightest sterile neutrino $N_1$ with mass in keV range can be a good candidate of warm dark matter. We found that the DW scenario where dark-matter $N_1$ is solely produced by the active-sterile transitions conflicts with the current cosmological constraints; X-ray and Ly-$\alpha$ constraints. To avoid this difficulty, we need some other production mechanism(s). Further, we
have shown that the oscillation between heavier sterile neutrinos $N_2$ and $N_3$ can be responsible to the baryon asymmetry of the universe. For the successful baryogenesis, their masses are almost degenerate and $M_{23} \sim 1$--10 GeV.

The $\nu$MSM predicts new fermions, three sterile neutrinos, which masses are within the reach of experimental energy. Thus, the primary issue is to verify these new states in various experiments and astrophysical observations$^{23}$.

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References


VII - Precision measurements
Precision Muon Physics Updates: $G_F$, $g_F$, and $g - 2$

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I report on a trio of precision muon experiments aimed at improvements in fundamental standard model parameters or in the search for new physics. New results from PSI yield an improved Fermi constant and the first unambiguous determination of the nucleon weak pseudoscalar coupling. These two experiments both involve precise measurements of the muon lifetime: $\tau_{\mu^+}$ obtains $G_F$, while $\tau_{\mu^-}$ in ultra-pure protium gas leads to $g_F$. The techniques and results are described. A planned next-generation muon anomalous magnetic moment experiment at BNL aims to significantly reduce the uncertainty in this test of new physics. The current status and a brief review of the physics case is given.

1 Introduction

The predictive power of the standard model depends on well-measured input parameters. In turn, tests of the standard model depend on comparisons of measurements to theoretical expectations. The difference is partially semantic; it can depend on what is known and when it is known. In the three vignettes presented below, we span the extremes. In the first chapter, a new precision muon lifetime measurement is described that updates the Fermi constant. Next, we report on a novel measurement of the $\mu p$ singlet capture rate, which is used to deduce the weak-nucleon pseudoscalar coupling; QCD-inspired theory predicts this quantity quite precisely. Finally, the muon anomaly is compared to the field-theory predictions that incorporate all known standard model processes. The difference between theory and experiment can be ascribed to new, missing, physics. In all three situations, the physics output implicitly depends on many "constants" and standard model parameters that are assumed to be known well enough, and the theory confidence is based on prior authenticity tests. We won’t dwell further on that subtlety here. The three sections which follow report a new $G_F$, a new $g_F$, and the physics reach to be expected in the LHC era from an improved $(g - 2)$ experiment.
2 The Fermi Constant

At this Electroweak Conference, we have repeatedly seen the Fermi constant, $G_F$, embedded in the predictions of the theory. That’s because $G_F$ is related to the strength of the weak interaction, much like the fine-structure constant $\alpha$ is to electromagnetism or “big G” is to gravity. With the assumption of universality in the weak interaction, any sufficiently precise weak measurement could be re-cast as a determination of $G_F$, as long as other parameters are well-enough established. By far, the most precise measure of $G_F$ is through ordinary muon decay, $\mu^+ \rightarrow e^+\nu_\mu\bar{\nu}_e$. The $G_F$ extracted from the lifetime (see below) is usually taken as the universal $G_F$ for any weak interaction rate. It is, in this sense, truly an “input” parameter so that the theory can predict less well measured quantities.

The last update of $G_F$ was made in 1984, despite vast improvements—in other EW measurements including the discovery of the weak bosons ($M_Z$ to 23 ppm!) and the top quark—since that time. The Fermi constant is related to the muon lifetime, $\tau_{\mu^+}$ by

$$ \frac{1}{\tau_{\mu^+}} = \frac{G_F^2 m_{\mu^+}^5}{192\pi^3} (1 + \Delta q), $$

where $\Delta q$ is the sum of phase space and both QED and hadronic radiative corrections. The relation

$$ \frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2} (1 + \Delta r) $$

connects $G_F$ to the electroweak gauge coupling, $g$. Here, $\Delta r$ represents the weak-boson-mediated tree-level and radiative corrections, which have been computed to second order.

The MuLan experiment uses an optimized “pulsed-muon-beam” technique to acquire the more than $10^{12}$ decays required for a 1 ppm measurement—the goal of the experiment. At PSI, the πE3 beamline is tuned to deliver $\sim 8 \times 10^8 \mu^+/s$ at $\sim 28$ MeV/c in dc operation. A customized time structure is applied to the beam using a fast kicker built for this experimental program. It works by applying a transverse electric field along a 1.5-m length of the beamline where the beam divergence is minimized. A 25 kV potential, switched on in 60 ns, is applied across two series-connected sets of aluminum plates to produce the field. When the plates are energized, the beam flux is reduced by a factor of more than 800. A 5 \( \mu \)s “accumulation period” $T_A$ is followed by a 22 \( \mu \)s “measuring period” $T_M$ (see Fig. 1). Muons are stopped in a thin ferromagnetic target, which has a high internal magnetic field to rapidly precess and decohere the initial polarization. At these rates and cycle, the experiment can accumulate muon decays at approximately 600 kHz. In 2006, $10^{12}$ decays were recorded, and in 2007 a similar data set will be collected using a different stopping target strategy.

The new result, alluded to in the introduction, is based on a much smaller sample of data obtained in 2004, when the experiment was being commissioned. A sample of $1.8 \times 10^{10}$ events, fit with a $\chi^2$/dof = 452.5/484 gives

$$ \tau_{\mu} (\text{MuLan}) = 2.197 \pm 0.013(21)(11) \mu \text{s} \quad (11.0 \text{ ppm}). $$

The first error is statistical and the second is systematic. The largest set of systematic uncertainties are absent for 2006 because of experimental changes in electronics and better performance of the kicker. However, challenging new systematic considerations enter at the precision of a 2 \( \mu \)s lifetime measurement. The 2006 data are presently being analyzed, and, just as in the 2004 case, the analysis is blinded by using a clock oscillator whose accurate and precise frequency is kept secret from the Collaboration until the analysis is complete. The new world average lifetime leads to an improved Fermi constant

$$ G_F = 1.166 \pm 0.006(5) \times 10^{-5} \text{ GeV}^{-2} \quad (5 \text{ ppm}). $$

3 The Weak Pseudoscalar Coupling

Muon capture on the proton—a weak interaction within a hadronic system—is a fundamental process whose rate is predicted following the symmetries of QCD. The process retains the familiar current-current $V - A$ form, but requires modified vector and axial vector pieces associated with form factors and having momentum transfer dependence. The rate for

$$\mu^- + p \rightarrow n + \nu_\mu$$

with the $\mu^- p$ atom in the singlet state is termed $\Lambda_S$. It is characterized by a matrix element where the vector and axial terms, assuming Lorentz and $T$-invariance (and no 2nd-class currents), can be written as

$$V^\alpha = \bar{u}_n (g_V q^2)^{\gamma^\alpha} + \frac{g_M(q^2)}{2m_n} q^{\beta\gamma} q_\beta) u_p$$

$$A^\alpha = \bar{u}_n (g_A q^2)^{\gamma^\alpha} \gamma_5 + \frac{g_P(q^2)}{m_n} q^{\beta} \gamma_5 u_p.$$  

Here $g_V, g_M, g_A$ are the vector, magnetic and axial-vector form factors, and $g_P$ is the induced pseudoscalar coupling, the least well known and the subject of a new measurement by the MuCap Collaboration.

The sensitivity of $\Lambda_S$ to these form factors is:

$$\frac{\delta \Lambda_S}{\Lambda_S} = 0.47 \frac{\delta g_V}{g_V} = 0.024\% \quad \frac{\delta \Lambda_S}{\Lambda_S} = 0.15 \frac{\delta g_M}{g_M} = 0.01\%$$

$$\frac{\delta \Lambda_S}{\Lambda_S} = 1.57 \frac{\delta g_A}{g_A} = 0.38\% \quad \frac{\delta \Lambda_S}{\Lambda_S} = 0.18 \frac{\delta g_P}{g_P} \approx 5\%$$

where it is clear that only $g_P$ is poorly known; a precision measurement of $\Lambda_S$ effectively deduces $g_P$.

On the theoretical side, $g_P$ is known to a few percent, giving $g_P(thy) = 8.26 \pm 0.23$. The experimental picture is murky and has been so for nearly 30 years. The reasons are associated with experimental complications from “muon chemistry.” The $\mu p$ system can form a $\mu p$ molecule in the ortho ($J = 1$, pp spins aligned) state, which can transition to the para ($J = 0$, pp spins anti-aligned) state.
pp spins anti-aligned) state at the rate $\lambda_{\text{op}}$. That means three muon-proton systems exist and their relative populations change with time. Further, the capture rate is different in all of them. In Fig. 2, $g_F$ at the fixed $q^2 = -0.88 ~ m_p^2$ of muon capture is given on the vertical axis. The heavy baryon chiral perturbation theory prediction is represented by the black band. The horizontal axis is $\lambda_{\text{op}}$, with two measurements and one theoretical calculation of $\lambda_{\text{op}}$ spanning the range $20 - 130 \times 10^3$ s$^{-1}$. Because of this large range, the ordinary muon capture (OMC) experiment performed in liquid hydrogen cannot be used to determine $g_F$—it is too sensitive to the unknown $\lambda_{\text{op}}$. The TRIUMF radiative muon capture (RMC) experiment is better, but its “high” value compared to theory has been controversial for many years. Is it an experimental issue or a challenge to theory? Could $\lambda_{\text{op}}$ be very high? The new MuCap result confirms the theory and is relatively immune to the uncertainty in $\lambda_{\text{op}}$.

The MuCap method uses a 10-bar, ultra-pure protium gas TPC target to image a muon stop in the gas vessel, well away from walls. The $\mu p p$ molecular formation is limited owing to the 1% density compared to liquid hydrogen. The TPC is surrounded by cylindrical wire chambers and a scintillating barrel hodoscope, which are used for tracking and timing of the decay electrons. A high-precision ($\sim 33$ ppm) negative muon lifetime experiment is performed, where the difference between $\tau_\mu$ and the free muon lifetime measured in MuLan is attributed to the capture rate, a difference of $\approx 0.15\%$. The rich technical challenges and solutions of MuCap are far too expansive to report here. The reader is urged to consult the PRI$^5$ for details.

The capture rate is $\lambda_S = 725.0 \pm 17.4$ s$^{-1}$. When systematics are accounted for and minor (known) corrections are applied, we find $g_F(\text{MuCap}) = 7.3 \pm 1.1$. Significantly more data have been obtained and the uncertainty will be halved in the future.

4 The Muon Anomaly

The muon anomaly is one of the most sensitive tests of the standard model because it can be measured and calculated precisely. Experiment and theory boast similar impressive uncertainties of $\approx 0.5$ ppm. When the numbers are compared, using the theoretical update given in the review of Miller, Roberts and deRafael$^{11}$, they do not agree, suggesting missing physics in the standard model evaluation or (let us hope not) a mistake in either of the numbers. The comparison gives

$$\Delta a_\mu^{(\text{today})} = a_\mu^{(\text{Exp})} - a_\mu^{(\text{SM})} = (29.5 \pm 8.8) \times 10^{-10}. \quad (10)$$
At this conference, Z. Zhang reviewed the theoretical contributions that make up the standard model value. These include QED, weak, and hadronic loops, with the latter carrying the largest uncertainty. Two main categories contribute nearly equally to the total theory uncertainty. The 1st-order hadronic vacuum polarization $a_{\mu}(HVP)$ is $690.1 \pm 4.7 \times 10^{-10}$. It comes from data—the absolute cross section for $e^+ e^- \rightarrow \text{hadrons}$, together with a well-known dispersion relation. Hadronic light-by-light (HLbL) is a second-order 4-point function that must be evaluated using a QCD-like model. The present best summary$^{12}$ of many efforts gives $a_{\mu}(\text{HLbL}) = (11.0 \pm 4.0) \times 10^{-10}$. Improvements—and even independent verification—of both numbers is important.

What could an anomalous moment having a $\sim 30 \times 10^{-10}$ departure from the standard model imply? Certainly, it points to new physics. Beyond that, it is only a part of the necessary clues required to “fingerprint” the source. Additional information will hail from direct measurements of masses and branching ratios at the LHC, from limits (or signals) from new charged lepton flavor violation experiments, from different EDM searches, and possibly other precision measurements, such as Möller scattering. Consider the landscape in the LHC era, by which time many of these projects will mature. For $(g - 2)$, a new experiment$^{13}$ E969 is approved at BNL but awaits funding. Its goal is a 2.5 or higher reduction in uncertainty. Improvements in theory are already on track with additional HVP data expected from radiative-return experiments at BaBar, KLOE and Belle. Some reduction in HLbL can be anticipated as well but here the path is less clear at the moment. For the sake of discussion, a $3.9 \times 10^{-10}$ uncertainty on the $\text{comparison}$ of experiment to theory is used as a future benchmark for new physics sensitivity.

In a recent White Paper$^{14}$, the physics case for such a scenario of improved precision is outlined. Let’s look at just one example$^{15}$ appropriate to this meeting—SUSY. Imagine a future in which the SPS1a reference point$^{16}$ is realized and the LHC has measured masses and a global fit has been performed to establish this model. Still, $\tan \beta$ will be largely unconstrained. In Fig. 3 a “blueband” plot is made with the reduction in $\chi^2$ versus $\tan \beta$ for the $a_{\mu}$ present (dark blue) and future (light blue) precisions considered when $a_{\mu}$ is included in the global fit. Compared to the LHC-alone limits (inside yellow bands), adding $a_{\mu}$ helps impressively. Other examples are given in the White Paper study.
5 Summary

In this brief report, we have announced two new measurements involving the muon lifetime. They update the Fermi constant and, for the first time, demonstrate unambiguous agreement between $g_\mu$ and fundamental QCD-inspired predictions. The new physics reach of the muon anomaly is already impressive with a 3.4 $\sigma$ significance on a deviation from the standard model. Plans for improved experiment and theory match nicely to the expected discoveries in the LHC era.

6 Acknowledgments

I was delighted to spend a week in the wonderful atmosphere of Moriond and I thank the organizers for the meeting and the invitation. I thank my colleagues on MuLan, MuCap and Muon $g-2$ for their creativity and scientific rigor, especially important qualities for high-precision measurements. Thanks to Lee Roberts and Steven Clayton for reading the draft. This work was sponsored, in part, by the U.S. National Science Foundation.

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13. A $(g-2)_{\mu}$ Experiment to $\pm 0.2$ ppm Precision, BNL Experiment E969, R.M. Carey, et al, (B.L. Roberts and D.W. Hertzog, co-spokespersons); See: http://g2pc1.bu.edu/roberts/Proposal969.pdf
16. SPS1a is a mSUGRA point with $m_0 = 100$ GeV; $m_{1/2} = 250$ GeV, $A_0 = -100$ and $\tan \beta = 10$. 
PRECISION MEASUREMENTS IN NEUTRON DECAY

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We present new precision measurements of angular correlation coefficients in polarized neutron decay. We have obtained values for the electron asymmetry coefficient $A$, the neutrino asymmetry coefficient $B$, and for the proton asymmetry coefficient $C$. In combination with other results, the new measurements are used to derive limits on "Physics beyond the Standard Model".

1 Introduction

Free neutrons decay with a mean lifetime $\tau_n$ of about 15 minutes in electron, proton, and electron anti-neutrino: $n \rightarrow e^- p \bar{\nu}_e$. The maximal kinetic energy of the electron is $E_{\text{max}}(e) = 782$ keV. For the proton, $E_{\text{max}}(p) = 780$ eV is three orders of magnitude smaller. Neutron decay – involving all particles of the first generation – constitutes an ideally suited laboratory to do particle physics at very low energies as it provides important information on the structure of the weak interaction and the underlying symmetries. It is a quite simple system where theoretical corrections are small and no nuclear structure effects have to be considered.

Our observables are angular correlations between neutron spin and the momentum of the three decay products. The decay probability $d\omega$ of polarized neutrons can be expressed in terms of two of these "asymmetries", the electron asymmetry $A$ and the neutrino asymmetry $B$

$$d\omega \propto |V_{ud}|^2 \left( 1 + a \frac{p_e p_{\nu}}{E_{\nu}} + \langle s_n \rangle \left[ A \frac{p_e}{E} + B \frac{p_{\nu}}{E_{\nu}} \right] \right).$$

(1)

$E, E_{\nu}, p_e, \text{ and } p_{\nu}$ are energy and momentum of electron and neutrino respectively, $\langle s_n \rangle$ denotes the neutron spin, and $|V_{ud}|$ is the first entry of the quark mixing matrix. $a$ is the correlation between electron and neutrino momentum. The proton asymmetry $C$ does not enter this expression, but it is kinematically coupled to $A$ and $B$ via $C = x_C (A + B)$, where $x_C = 0.27484$ is a kinematical factor. In the standard $V - A$ formulation of weak interactions, all correlations
are functions of one single parameter $\lambda = g_A/g_V$, the ratio of axial-vector and vector coupling constant (assuming $\lambda$ to be real and neglecting weak magnetism and other recoil effects):

$$a = \frac{1 - \lambda^2}{1 + 3\lambda^2}, \quad A = -2\frac{\lambda^2 + \lambda}{1 + 3\lambda^2}, \quad B = 2\frac{\lambda^2 - \lambda}{1 + 3\lambda^2}, \quad C = xC\frac{4\lambda}{1 + 3\lambda^2}. \quad (2)$$

The precise determination of $\lambda$ is important as the calculation of many processes in cosmology (e.g. primordial element formation), astronomy (e.g. solar cycle, neutron star formation), and particle physics (e.g. neutrino detectors, neutrino scattering) depends on this parameter. The observable most sensitive to $\lambda$ is the electron asymmetry $A$, however, it can be also extracted from measurements of $a$ and $C$.

Within the Standard Model, neutron decay can be described with the two parameters $\lambda$ and $V_{ud}$ only. But since much more observables are accessible (various correlation coefficients and the lifetime $\tau_n$) the problem is overdetermined, and neutron decay can be used to test the Standard Model and to search for new physics. In section 3, we will present neutron limits for additional right-handed $(V + A)$ currents and anomalous couplings (scalar and tensor) in the interaction.

2 Experiment and Results

The measurements were performed using the electron spectrometer PERKEO II. It features a pair of superconducting coils in a split pair configuration that generate a slightly decreasing magnetic field ($B_{\text{max}} \approx 1$ T) perpendicular to the spin polarized neutron beam crossing the spectrometer (cf. fig. 1). The instrument was installed at the cold neutron beam position PF1B at the Institut Laue-Langevin (ILL), Grenoble.

The magnetic field fulfills several functions: It separates the full solid angle into two hemispheres since the spins align with the field lines: The momentum projection of the decay particle onto the neutron spin determines whether the particle is emitted in neutron spin direction or against it. The field guides the charged decay products onto the two detectors, installed at both sides next to the beam, providing full $2 \times 2 \pi$ detection. The detectors consist of plastic scintillators with photomultiplier readout. In order to measure the very low energetic protons with the same detector (necessary for neutrino and proton asymmetry), we developed a special

![Figure 1: Electron spectrometer PERKEO II: Transversally polarized neutrons transit the instrument, their spins are aligned with the magnetic field, which separates the full solid angle into two hemispheres covered by a detector each.](image-url)
setup to convert protons into electrons\textsuperscript{3,4}: Protons were accelerated onto a very thin carbon foil on negative high voltage where they had enough ionization power to generate one or more secondary electrons. These could be detected by the scintillators.

2.1 Electron Asymmetry $A$ and $\lambda$

The experimental signature of the electron asymmetry is

$$A_{\text{exp}}(E) = \frac{N^+(E) - N^-(E)}{N^+(E) + N^-(E)} = \frac{1}{2} \frac{v}{c} A P F,$$

where $N(E)$ is the number of electrons with energy $E$ and the sign denotes whether the electron was emitted in $(+)$ or against $(-)$ neutron spin direction. This expression holds for both detectors separately since a spinflipper was used to periodically turn the neutron spin by $180^\circ$. It is related to the angular correlation coefficient $A$ via the second equation in (3). $v/c$ is the electron velocity in terms of the speed of light, $P$ the neutron polarization, and $F$ the spinflipper efficiency. With new methods to polarize the beam and to analyze the beam polarization\textsuperscript{5} we managed to achieve a much improved beam polarization compared to former measurements.

Background generated in the neutron collimation system and at the beamstop at the end of the installation was heavily suppressed. The remaining background leads to a small uncertainty of 0.1 %. For future measurements, the spectrometer PERKEO III\textsuperscript{6} has been developed: It can be operated with a chopped neutron beam allowing to acquire data only when no background is generated.

About 160 million events were recorded and the measurement is still limited by statistics. Details on the experiment can be found in\textsuperscript{7}. For the following analysis, we will use the value $A_{\Pi} = -0.1193(4)$ - corresponding to $\lambda_{\Pi} = -1.2749(11)$, the average of the new preliminary result and former PERKEO II measurements\textsuperscript{8}.

2.2 Neutrino Asymmetry $B$

In our setup, the neutrino could not be detected directly, thus it had to be reconstructed from a coincident measurement of electron and proton. The case where both are emitted into the same hemisphere is most sensitive to the neutrino asymmetry $B^3$. Here, the experimental asymmetry is defined using the electron spectra $Q^j$

$$B_{\text{exp}}(E) = \frac{Q^--(E) - Q^{++}(E)}{Q^--(E) + Q^{++}(E)},$$

where the first sign indicates the emission direction of the electron, the second denotes the proton.

A fit to the combined data is shown in figure 2. The result (for details cf.\textsuperscript{4})

$$B_{\Pi} = 0.9802(50)$$

is almost independent from detector calibration due to the flat characteristics of the spectrum. It has an uncertainty that is comparable to the most precise measurement so far\textsuperscript{10} but has significantly lower corrections. The result is limited by statistics and the error in the relative position between the magnetic field of the spectrometer and the neutron beam maximum. A misalignment would cause some charged particles to be reflected at the increasing magnetic field ("magnetic mirror") leading to signals in the wrong detector.

The new result agrees with the Standard Model expectation (calculated with eq. (2) and the current PDG value for $\lambda$) and previous measurements\textsuperscript{10,11}. Averaging all yields the new world mean value with an uncertainty lowered by 25 %:

$$B_{\text{mean}} = 0.9807(30).$$
2.3 Proton Asymmetry C

The combined electron-proton detector only allows to measure the emission direction of the proton, a determination of its energy is impossible since it is much smaller than the electron energy. Therefore we have to obtain the proton asymmetry $C$ from the coincident measurement in an integral way. For both detectors, altogether four electron energy spectra generated with certain conditions on the emission direction of electrons (first) and protons (second sign) are available:

$$Q^{++}(E), \quad Q^{-+}(E), \quad Q^{--}(E), \quad Q^{+-}(E). \quad (7)$$

Now, the proton asymmetry is defined using the integrals of these spectra

$$C = \frac{\int (Q^{++}(E) + Q^{-+}(E)) \, dE - \int (Q^{--}(E) + Q^{+-}(E)) \, dE}{\int (Q^{++}(E) + Q^{--}(E)) \, dE + \int (Q^{-+}(E) + Q^{+-}(E)) \, dE} \quad (8)$$

The integrals were obtained from one-parameter fits to the $Q^{ij}$ spectra in a fit region that was well above the detection threshold and remaining background contributions. After extrapolation to the whole energy range, the functions were integrated. This yields the result $^{12}$

$$C_{_{\Pi}} = -0.2377(36). \quad (9)$$

With an uncertainty of 1.5 %, this constitutes the first precision measurement of this observable. The error is dominated by energy calibration and extrapolation, however, at the moment the proton asymmetry is known more precisely than the electron-neutrino correlation $\alpha$.

3 Neutron Decay Limits on New Physics

We use the new results in combination with other measurements to set limits on possible contributions of new physics to neutron decay. Another important input parameter is the neutron lifetime $\tau_n$. At the moment, however, the experimental situation concerning $\tau_n$ is quite unclear, since there is one new measurement $^{13}$ deviating by 6.5 $\sigma$ from the average of previous results. In order to account for this discrepancy we enlarge all quoted errors with a scaling factor 2.5 to obtain statistical agreement ($\chi^2/NDF = 1$) and get the average $\tau_{\text{mean}} = 882.0(14)$ s.

3.1 Right-Handed Currents

Parity is maximally violated in the Standard Model, i.e. the weak interaction couples only to left-handed particles. However, the model only describes but gives no intrinsic motivation
for parity violation. According to simple extensions of the Standard Model, the Left-Right Symmetric Models\textsuperscript{14,15}, parity violation stems from a spontaneous symmetry breaking at mass scale $m_2$. Below this energy, the interaction is mediated by the usual $W_L$-bosons, however, there should be additional heavy bosons $W_R$, remnants of a right-handed SU(2)$_R$ group. The Manifest Left-Right Symmetric Model\textsuperscript{16} assumes that left- and right-handed quark mixing matrices are equal. Here, the weak eigenstates $W_{L,R}$ are linear combinations of the mass eigenstates $W_{1,2}$,

$$W_L = \cos \zeta \ W_1 - \sin \zeta \ W_2 \quad \text{and} \quad W_R = e^{i\phi} \sin \zeta \ W_1 + e^{i\phi} \cos \zeta \ W_2,$$

with mixing angle $\zeta$ and a CP violating phase $\phi$ (we will neglect $\phi$ in the following since it has no observable effect). Additional parameters of the theory are the mass ratio $\delta \equiv m_2^2/m_1^2$ and the coupling constant ratio $\lambda' = g'_{4}/g'_{\nu}$.

If one extends equations (2) and the expression for the lifetime $\tau_n$ to account for possible right-handed admixtures, one can generate exclusion plots and derive limits. The neutron decay result for the input parameters $A_{PL}$, $B_{\text{mean}}$, and $\tau_{\text{mean}}$ is given in figure 3 showing a projection along the $\lambda'$ axis onto the $\zeta - \delta$ plane. The 90 % confidence level (CL) limits are $-0.1968 < \zeta < 0.0040$ and $\delta < 0.0885$ which yields $m_2 > 270 \text{ GeV}$.

There are tighter limits on the mass $m_2$ from direct $W'$ searches at Tevatron\textsuperscript{17}, and better constraints on $\zeta$ from muon decay tests\textsuperscript{18}, however, in the mass range not excluded by the collider results the neutron limits for the mixing angle $\zeta$ are more stringent. And when one considers more general right-left symmetric models, results from $\beta$-decay, muon decay, and direct searches are complementary\textsuperscript{19}.

\subsection{3.2 Scalar and Tensor Couplings}

In the framework of the Standard Model, the weak interaction includes only vector ($V$) and axial-vector ($A$) couplings. However, the most general Lorentz invariant Lagrangian\textsuperscript{20}

$$\mathcal{L} = \sum_k (\not\mathcal{P}_k \not\mathcal{V}_k \not\nu) \left( \not\mathcal{O}_k (g_k + g'_k \gamma^5) \nu \right) + \text{h.c.}$$

allows also scalar ($S$) and tensor ($T$) contributions, where the operator $\mathcal{O}_k$ ($k = V, A, S, T$) describes the kind of interaction. Again, equations (2) were extended to account for additional

\begin{figure}[h]
\begin{center}
\includegraphics[width=0.4\textwidth]{exclusion_plot}
\caption{Exclusion plot on possible admixtures of right-handed ($V + A$) couplings in the weak interaction derived from neutron decay observables. The Standard Model, $\zeta = \delta = 0$, is included with 95 \% CL.}
\end{center}
\end{figure}

\begin{figure}[h]
\begin{center}
\includegraphics[width=0.4\textwidth]{neutron_decay_limits}
\caption{Neutron decay limits on anomalous (scalar $g_S$ and tensor $g_T$) couplings in the weak Lagrangian. The Standard Model $g_S = g_T = 0$ is included in the 90 \% CL contour.}
\end{center}
\end{figure}
couplings $g_S, g_T$ in the simple Right-Handed Scalar and Tensor Model\textsuperscript{10}. It assumes $g'_V/g_V = 1$, $g'_A/g_A = 1, g'_S/g_V = -g_S/g_V$, and $g'_T/g_A = -g_T/g_A$ leading to left-handed $V, A$ and right-handed $S, T$ couplings. The result of a $\chi^2$-scan based on the input parameters $A_{\text{pp}}, B_{\text{pp}}, C_{\text{pp}}, a = -0.103(4)\text{, and } \tau_{\text{mean}}$ is shown in fig. 4. The 90 \% CL limits, $|g_S/g_V| < 0.130$ and $|g_T/g_A| < 0.0948$, agree with the Standard Model case $g_S = g_T = 0$.

Scalar contributions seem to be almost excluded by an analysis of superallowed $\beta$-decays giving the constraint $|g_S/g_V| < 0.0013$ when the conserved vector current hypothesis (CVC) is assumed\textsuperscript{22}. For tensor contributions, however, neutron decay provides excellent limits.

Acknowledgments

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The current status of the experimental measurements and theoretical predictions of the anomalous magnetic moment of the muon $a_\mu$ is briefly reviewed. The emphasis is put on the evaluation of the hadronic contribution to $a_\mu$ as it has the largest uncertainty among all Standard Model contributions. The precision of the hadronic contribution is driven by the input $e^+e^-$ data predominantly from the $\pi^+\pi^-$ channel. Including the latest experimental data on $e^+e^-$ annihilation into hadrons from CMD2 and SND for the $\pi^+\pi^-$ channel and BaBar for multihadron final states, the updated Standard Model prediction disagrees with the measurement dominated by BNL by 3.3 standard deviations, with the theoretical precision exceeding the experimental one.

1 Introduction

For a charged elementary particle with $1/2$ intrinsic spin such as muon, its magnetic dipole moment $\vec{\mu}$ is aligned with its spin $\vec{s}$ as:

$$\vec{\mu} = g \left( \frac{q}{2m} \right) \vec{s},$$

(1)

where $q = \pm e$ is the charge of the particle in unit of the electron charge and $g$ is the gyromagnetic ratio. In the classic Dirac theory, $g = 2$. In the Standard Model (SM), quantum loop effects induce a small correction, which is quantified by $a_\mu = (g_\mu - 2)/2$, the so-called anomalous magnetic moment or the magnetic anomaly.

There has been a long history in measuring and calculating $a_\mu$. In particular the steadily improving precision of both the measurements and the predictions of $a_\mu$ and the disagreement observed between the two have led the study of $a_\mu$ one of the most active research fields in particle physics in recent years.
The paper is organized as follows. In section 2, the measurement history and the current world average value of $\mu_p$ are presented. In section 3, different components of the SM contributions to $\mu_p$ are reviewed. Section 4 is reserved for discussions followed by conclusion and prospects in section 5.

2 The Measurement of $\mu_p$

A compilation of the major experimental efforts in measuring $\mu_p$ over the last five decades is given in Table 1 (a modified version of Table 1 from a recent review article \textsuperscript{1}). Starting from the experiment at the Columbia-Nevis cyclotron, where the spin rotation of a muon in a magnetic field was observed for the first time, the experimental precision of $\mu_p$ has seen constant improvement first through three experiments at CERN in the sixties and seventies and more recently with ES21 at the Brookhaven National Laboratory (BNL). The current world average value reaches a relative precision of 0.54 ppm (parts per million).

Table 1: Measurements of the muon magnetic anomaly $\mu_p$, where the value in parentheses stands for either the total experimental error or the statistical and systematic ones.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>Beam</th>
<th>Measurement</th>
<th>$\delta \mu_p/\mu_p$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Columbia-Nevis(1957) \textsuperscript{2}</td>
<td>$\mu^+$</td>
<td>$g = 2.00 \pm 0.10$</td>
<td>0.001 13$^{+16}_{-12}$</td>
</tr>
<tr>
<td>Columbia-Nevis(1959) \textsuperscript{3}</td>
<td>$\mu^+$</td>
<td>0.001 145(22)</td>
<td>1.9%</td>
</tr>
<tr>
<td>CERN 1(1961) \textsuperscript{4}</td>
<td>$\mu^+$</td>
<td>0.001 162(5)</td>
<td>0.43%</td>
</tr>
<tr>
<td>CERN 1(1962) \textsuperscript{5}</td>
<td>$\mu^+$</td>
<td>0.001 166 16(31)</td>
<td>265 ppm</td>
</tr>
<tr>
<td>CERN 3(1968) \textsuperscript{6}</td>
<td>$\mu^+$</td>
<td>0.001 165 896(27)</td>
<td>23 ppm</td>
</tr>
<tr>
<td>CERN 3(1975) \textsuperscript{7}</td>
<td>$\mu^+$</td>
<td>0.001 165 911(11)</td>
<td>7.3 ppm</td>
</tr>
<tr>
<td>CERN 3(1979) \textsuperscript{8}</td>
<td>$\mu^+$</td>
<td>0.001 165 919 1(59)</td>
<td>5 ppm</td>
</tr>
<tr>
<td>BNL ES21(2000) \textsuperscript{9}</td>
<td>$\mu^+$</td>
<td>0.001 165 920 2(16)</td>
<td>1.3 ppm</td>
</tr>
<tr>
<td>BNL ES21(2001) \textsuperscript{10}</td>
<td>$\mu^+$</td>
<td>0.001 165 920 3(8)</td>
<td>0.7 ppm</td>
</tr>
<tr>
<td>BNL ES21(2002) \textsuperscript{11}</td>
<td>$\mu^+$</td>
<td>0.001 165 921 4(8)</td>
<td>0.7 ppm</td>
</tr>
<tr>
<td>BNL ES21(2004) \textsuperscript{12}</td>
<td>$\mu^+$</td>
<td>0.001 165 920 80(63)</td>
<td>0.54 ppm</td>
</tr>
</tbody>
</table>

The muon magnetic anomaly $\mu_p$ in all modern experiments is determined by the following method. For an ensemble of polarized muons which are moving in a storage ring in a highly uniform magnetic field $\vec{B}$ (perpendicular to muon spin and orbit plane) and a vertically focusing quadrupole field $\vec{E}$, the frequency difference $\omega_\alpha$ between the spin precession $\omega_\alpha$ and the cyclotron motion $\omega_\gamma$ is described by

$$\omega_\alpha \equiv \omega_\gamma - \omega_\alpha = \frac{e}{m_p c} \left[ \mu_p \vec{B} - \left( \frac{1}{\gamma - 1} \right) (\vec{\beta} \times \vec{E}) \right], \quad \text{when} \quad \vec{B} \cdot \vec{\beta} = \vec{E} \cdot \vec{\beta} = 0 \quad (2)$$

where $\vec{\beta}$ represents the muon direction. The second term in parentheses vanishes at $\gamma = 29.3$ (the magic momentum) and the electrostatic focusing does not affect the spin. The key to the experiment is to determine frequency $\omega_\alpha$ to high precision and to measure the average magnetic field to equal or better precision.

In comparison with the electron magnetic anomaly, $\alpha_e$ is more precisely measured \textsuperscript{14} (0.7 ppb), but $\mu_p$ is more sensitive to new physics effects by about $m^2_{\mu}/m^2_e \approx 40 000$ because of its large mass value.
3 Prediction of the Standard Model Contributions

In the SM, the muon magnetic anomaly $a_\mu$ receives contributions from all electromagnetic (QED), weak and strong (hadronic) sectors and can be conveniently written as:

$$a^{\text{SM}}_\mu = a^{\text{QED}}_\mu + a^{\text{weak}}_\mu + a^{\text{had}}_\mu.$$  \hspace{1cm} (3)

Their representative diagrams are shown in Fig. 1, which also includes two example contributions from new particles in supersymmetry models. Thus comparison of the precision measurement and theory tests the validity of the SM at its quantum loop level and probes effects of new physics.

3.1 QED and Weak Contributions

The QED correction, which includes all photonic and leptonic ($e$, $\mu$ and $\tau$) loops, is by far the dominant contribution in the SM:

$$a^{\text{QED}}_\mu = \frac{\alpha}{2\pi} + 0.765857410(27) \left( \frac{\alpha}{\pi} \right)^2 + 24.05050964(43) \left( \frac{\alpha}{\pi} \right)^3 + 130.9916(80) \left( \frac{\alpha}{\pi} \right)^4 + $$

$$+ 663(20) \left( \frac{\alpha}{\pi} \right)^5 + \cdots,$$ \hspace{1cm} (4)

where the lowest-order Schwinger term ($\alpha/2\pi$) was known since 1948 $^{14}$, the coefficients are analytically known for terms up to ($\alpha/\pi$)$^3$, numerically calculated for the fourth term and recently estimated for the fifth term $^{15}$. Using $\alpha$ extracted from the latest $a_e$ measurement $^{14}$, one has

$$a^{\text{QED}}_\mu = 116584718.09(0.14)_{\text{exp}}(0.08)_{\delta \alpha} \times 10^{-11}.$$  \hspace{1cm} (5)

The weak contributions, involving heavy $Z$, $W^\pm$ or Higgs particles, are suppressed by at least a factor $\frac{\alpha}{\pi} \frac{m^2}{M_W^2} \simeq 4 \times 10^{-9}$. At one-loop order,

$$a^{\text{weak}}_{\mu} \ [\text{1-loop}] = \frac{G_F m^2}{8 \sqrt{2} \pi^2} \left[ \frac{5}{3} + \frac{1}{3} (1 - 4 \sin^2 \theta_W)^2 + O \left( \frac{m_H^2}{M_W^2} \right) + O \left( \frac{m^2}{M_H^2} \right) \right]$$ \hspace{1cm} (6)
\[ = 194.8 \times 10^{-11}, \quad \text{for } \sin^2\theta_W \equiv 1 - \frac{M_W^2}{M_Z^2} = 0.223. \quad (7) \]

Two-loop corrections are relatively large and negative
\[ \alpha_{\mu}^{\text{weak}[2\text{-loop}]} = -40.7(1.0)(1.8) \times 10^{-11}, \quad (8) \]
where the errors stem from quark triangle loops and the assumed Higgs mass range \( M_H = 150^{+100}_{-40} \) GeV. The three-loop leading logarithms are negligible, \( \mathcal{O}(10^{-12}) \), implying in total
\[ \alpha_{\mu}^{\text{weak}} = 154(1)(2) \times 10^{-11}. \quad (9) \]

### 3.2 Hadronic Contributions

The hadronic contributions are associated with quark and gluon loops. They cannot be calculated from first principles because of the low energy scale involved. Fortunately, owing to unitarity and to the analyticity of the vacuum polarization function, the lowest-order hadronic vacuum polarization contribution to \( \alpha_{\mu} \) can be computed via the dispersion integral\(^{17}\) using the ratio \( R^{(3)}(s) \) of the bare cross section\(^{d}\) for \( e^+e^- \) annihilation into hadrons to the pointlike muon pair cross section at center-of-mass energy \( \sqrt{s} \)
\[ \alpha_{\mu}^{\text{had,LO}} = \frac{1}{3} \left( \frac{\alpha}{\pi} \right)^2 \int_{m_\mu^2}^\infty ds \frac{K(s)}{s} R^{(3)}(s), \quad (10) \]
where \( K(s) \) is the QED kernel\(^{18}\) \( K(s) = x^2 \left[ 1 - \frac{x^2}{2} \right] + (1 + x)^2 \left[ 1 + \frac{1}{2}x \right] \left[ \ln(1+x) - x + \frac{x^2}{2} \right] + x^2 \ln\frac{1 - \beta_\mu}{1 + \beta_\mu} \), with \( x = \frac{1 - \beta_\mu}{1 + \beta_\mu} \) and \( \beta_\mu = \left( 1 - \frac{4m_e^2}{s} \right)^{1/2} \). The kernel function \( K(s) \sim \frac{1}{s} \) gives weight to the low energy part of the integral. About 91\% of the total contribution to \( \alpha_{\mu}^{\text{had,LO}} \) is accumulated at \( \sqrt{s} \) below 1.8 GeV and 73\% of \( \alpha_{\mu}^{\text{had,LO}} \) is covered by the \( \pi\pi \) final state, which is dominated by the \( \rho(770) \) resonance.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>( N_{\text{data}} )</th>
<th>Energy range (GeV)</th>
<th>( \delta(\text{stat.}) )</th>
<th>( \delta(\text{syst.}) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>DM1 (1978) (^{22})</td>
<td>16</td>
<td>0.483 - 1.096</td>
<td>(6.6 - 40)%</td>
<td>2.2%</td>
</tr>
<tr>
<td>TOF (1981) (^{23})</td>
<td>4</td>
<td>0.400 - 0.460</td>
<td>(14 - 20)%</td>
<td>5%</td>
</tr>
<tr>
<td>OLYA (1979, 1985) (^{24,25})</td>
<td>2 + 77</td>
<td>0.400 - 1.397</td>
<td>(2.3 - 35)%</td>
<td>4%</td>
</tr>
<tr>
<td>CMD (1985) (^{26})</td>
<td>24</td>
<td>0.360 - 0.820</td>
<td>(4.1 - 10.8)%</td>
<td>2%</td>
</tr>
<tr>
<td>DM2 (1989) (^{26})</td>
<td>17</td>
<td>1.350 - 2.215</td>
<td>(17.6 - 100)%</td>
<td>12%</td>
</tr>
<tr>
<td>CMD2 (2003) (^{27})</td>
<td>43</td>
<td>0.611 - 0.962</td>
<td>(1.8 - 14.1)%</td>
<td>0.6%</td>
</tr>
<tr>
<td>KLOE (2005) (^{28})</td>
<td>60</td>
<td>0.600 - 0.970</td>
<td>(0.5 - 2.1)%</td>
<td>(1.2 - 3.8)%</td>
</tr>
<tr>
<td>SND (2006) (^{29})</td>
<td>45</td>
<td>0.390 - 0.970</td>
<td>(0.5 - 2.1)%</td>
<td>(1.2 - 3.8)%</td>
</tr>
<tr>
<td>CMD2(_{\text{low}}) (2006) (^{30})</td>
<td>10</td>
<td>0.370 - 0.520</td>
<td>(4.5 - 7)%</td>
<td>0.7%</td>
</tr>
<tr>
<td>CMD2(_{\text{rho}}) (2006) (^{31})</td>
<td>29</td>
<td>0.600 - 0.970</td>
<td>(0.5 - 4.1)%</td>
<td>0.8%</td>
</tr>
<tr>
<td>CMD2(_{\text{high}}) (2006) (^{32})</td>
<td>36</td>
<td>0.980 - 1.380</td>
<td>(4.5 - 18.4)%</td>
<td>(1.2 - 4.2)%</td>
</tr>
</tbody>
</table>

A detailed compilation of all the experimental data used in the evaluation of the dispersion integral prior to 2004 is provided in Refs.\(^{20,21}\). Since then, a few precise measurements have been published. A list of experiments for the dominant \( \pi\pi \) channel is shown in Table 2.

The \( \pi\pi \) data are compared in Fig. 2. Closer inspections show that the most precise measure-

\(^{a}\)The bare cross section is defined as the measured cross section corrected for initial state radiation, electron vertex contributions and vacuum polarization effects in the photon propagator but with photon radiation in the final state included\(^{19}\).
Figure 2: Comparison of $\pi^+\pi^-$ spectral functions expressed as $e^+e^-$ cross sections. The band corresponds to combined data used in the numerical integration.

Measurements from the annihilation experiments SND and CMD2 at Novosibirsk are in good agreement. They differ however in shape with those measured by KLOE using the radiative return method at DAΦNE\(^{28}\) (see Sec. 4.1). Before this is clarified, the KLOE data are not used in some of the recent evaluations of $a_{\mu}^{\text{had},\text{LO}}$.

In addition to the dominant $\pi\pi$ mode, results from the B\(\Lambda\)B\(\Phi\)/B\(\Phi\)R experiments are being produced on multihadron final states using also radiative return\(^{33}\). Benefiting from its big initial center-of-mass energy of 10.6 GeV, hard-radiated photon detected at large angle and high statistics data sample, the B\(\Lambda\)B\(\Phi\)/B\(\Phi\)R measurements are precise over the whole mass range. One example is shown in Fig. 3 in comparison with earlier measurements.

Including these new input $e^+e^-$ data, a preliminary update\(^{b}\) of $a_{\mu}^{\text{had},\text{LO}}$ is performed\(^{40}\) and shown in Table 3. There is no new tau data since the previous evaluation, therefore the $\tau$ based calculation is taken directly from Ref.\(^{21}\). The evaluation using $\tau$ data is made\(^{34}\) by relating the vector spectral functions from $\tau \rightarrow \nu_\tau \tau$ hadrons decays to isovector $e^+e^- \rightarrow$ hadrons cross sections by isospin rotation. All known isospin breaking effects are then taken into account\(^{20,21}\).

The higher order (NLO) hadronic contributions $a_{\mu}^{\text{had},\text{NLO}}$ involve one hadronic vacuum polarization insertion with an additional loop (either photonic or leptonic or another hadronic vacuum polarization). They can be evaluated\(^{35}\) with the same $e^+e^- \rightarrow$ hadrons data sets used for $a_{\mu}^{\text{had},\text{LO}}$. The numerical value\(^{36}\) reads

$$a_{\mu}^{\text{had},\text{NLO}} = -9.79(0.09)_{\exp}(0.03)_{\text{rad}} \times 10^{-10} \tag{11}$$

where the first and second errors correspond respectively to the experimental uncertainty of the $e^+e^-$ data and the radiative correction uncertainty.

Another higher order hadronic contribution to $a_{\mu}$ is from the hadronic light-by-light scattering (illustrated with the lower figure in the second column in Fig. 1). Since it invokes a four-point

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\(^{b}\)It is preliminary as some of the new $e^+e^-$ data used were still in their preliminary form.
correlation function, a dispersion relation approach using data is not possible at present. Instead, calculations involving pole insertions (or Goldstone boson exchanges), short distance quark loops and charged pion (and kaon) loops have been individually performed in a large $N_c$ QCD approach. A representative value used in Ref. $^20$ was $a_{\mu}^{had,LBL} = 8.6(3.5) \times 10^{-10}$.

A new analysis $^{38}$, which takes into account the proper matching of asymptotic short-distance behavior of pseudoscalar and axial-vector contributions with the free quark loop behavior, leads to $a_{\mu}^{had,LBL} = 13.6(2.5) \times 10^{-10}$. However, as pointed out in Ref. $^{30}$, several small but negative contributions such as charged pion loops and scalar resonances were not included in the latter calculation, thus in a recent update evaluation $^{40}$ of $a_{\mu}$, the following value

$$a_{\mu}^{had,LBL} = 12.0(3.5) \times 10^{-10}$$

was used. This is consistent with the value $a_{\mu}^{had,LBL} = 11(4) \times 10^{-11}$, suggested in Ref. $^{41}$. The uncertainty $a_{\mu}^{had,LBL}$, being the second largest one next to $a_{\mu}^{had,LO}$, clearly needs improvement in the near future.

$^5$A different evaluation $^{44}$ used directly the value of Ref. $^{38}$ of $13.6(2.5) \times 10^{-10}$. 

![Graph](image_url)
Adding all SM contributions together, the comparison from recent evaluations with the measurement is shown in Fig. 4. While the $\tau$ data-based calculation agrees with the measurement within the errors, the $e^+e^-$ data-based evaluations show a deviation of around 3.3 standard deviations.

![Figure 4: Comparison of recent theoretical evaluations of $a_\mu$ with the BNL measurement.](image)

4 Discussions

4.1 Tau Data versus $e^+e^-$ Data

The $\tau$ data used in the $a_\mu$ evaluation is the averaged one from the LEP experiments ALEPH and OPAL, and the CLEO experiment. The data are compared in Ref. and found in good agreement in particular for the two most precise data from ALEPH and CLEO. These data are complementary as the ALEPH data are more precise below the $\rho$ peak while CLEO has the better precision above.

A comparison between the averaged $\tau$ data and the $e^+e^-$ for the dominant $\pi\pi$ mode is shown in Fig. 5. The difference of $5 - 10\%$ in the energy region of $0.65 - 1.0$ GeV$^2$ is clearly visible. The difference with KLOE is even more pronounced.

4.2 CVC

An alternative way of comparing $\tau$ and $e^+e^-$ data is to compare measurements of branching fractions $B$ in $\tau$ decays with their expectations from CVC (Conserved Vector Current) using $e^+e^-$ spectral functions, duly corrected for isospin breaking effects. The advantage of a such comparison is that the measurements of $B$ are more robust than the spectral functions as the latter ones depend on the experimental resolution and require a numerically delicate unfolding.

The comparison for $\pi\pi$ mode revealing a discrepancy of 4.5 standard deviations is shown in Fig. 6. Similar comparisons for decay modes $\tau^- \rightarrow \nu_\tau \pi^- 3\pi^0$ and $\tau^- \rightarrow \nu_\tau 2\pi^- \pi^+ \pi^0$ have also been made and the differences with the corresponding $e^+e^-$ data are found respectively at 0.7 and 3.6 standard deviations.
5 Conclusion and Prospects

The muon magnetic anomaly $a_{\mu}$ is one of the most precisely known quantities both experimentally and theoretically in the SM. Incorporating new $e^+e^-$ data from CMD2 and SND for $\pi\pi$ mode and from BABAR for multihadronic modes, new SM determinations of $a_{\mu}$ have been obtained with a theoretical precision exceeding for the first time in recent years the experimental one. The SM prediction is found to be smaller than the measurement by about 3.3 standard deviations. Unfortunately one can not draw a definitive conclusion for the moment as the $\tau$ data based prediction is in agreement with the measurement.

Therefore it is extremely important that one clarifies the discrepancy between the $e^+e^-$ and $\tau$ data in particular on the $\pi\pi$ mode. There are a number of possibilities: (1) (the normalization of) the $e^+e^-$ data is wrong, (2) the tau data are wrong, (3) both are correct but there are unaccounted effects which explain the discrepancy between the two.

Possibility (1) may be also related to the current difference (mainly on the shape of the
spectral functions) between CMD2/SND data and KLOE data obtained respectively from the beam energy scan method and the radiative return method. This difference is expected to be resolved soon as KLOE has more and high quality data to be analyzed. In addition, reduced systematics uncertainties can be achieved if the measurement is made by normalizing the $\pi\pi$ data to $\mu\mu$ instead of to luminosity using the large angle Bhabha process.

The long awaited high precision measurement in $\pi\pi$ mode from BABAR using also the radiative return method will certainly help in clarifying some of the issues.

On the tau side, further improvement on the high mass part of the spectral functions is expected from large statistical data samples available at $B$ factories and a $\tau$-charm factory.

While the leading hadronic uncertainty gets improved with the forthcoming high precision $e^+e^-$ (and $\tau$) data, the next item awaiting for significant improvement concerns the uncertainty on the light-by-light scattering contribution $a^{\text{had},\text{LBL}}$.

Given the fact that the theoretical error is already smaller than the experimental one, it is timely to improve the latter. Indeed there is a new project BNL-E969 allowing to reduce the current error by more than a factor two down to 0.24 ppm. We are looking forward that the project gets funded very soon.

Acknowledgments

The author wishes to thank the organizers of the conference for the invitation and M. Davier, S. Eidelman and A. Höcker for the fruitful collaboration.

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Measurement of EPR-type flavour entanglement in $\Upsilon(4S) \rightarrow B^0 \bar{B}^0$ decays

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The neutral $B$-meson pair produced at the $\Upsilon(4S)$ should exhibit a non-local correlation of the type discussed by Einstein, Podolski, and Rosen. The time-dependent flavour asymmetry of the $B$ mesons decaying into flavour eigenstates will be used to test such a correlation. The asymmetry obtained from semileptonic $B^0$ decays is in agreement with the prediction from quantum mechanics and far away from the predictions of local realism models. We also test for possible partial decoherence effects. Our results are consistent with no decoherence.

1 Introduction

The concept of entangled states (i.e., states which cannot be represented as product states of their parts) was born in the '30 in the midst of several conceptual difficulties with Quantum Mechanics (QM). In 1935 Einstein, Podolski, and Rosen (EPR) arrived at the conclusion that QM could not be a “complete” theory. EPR considered a pair of particles produced by the same interaction, subsequently freely propagating in space but still linked by momentum conservation. EPR found a contradiction when realism and locality are applied to the predictions of QM on a couple of non-commuting observables (position and momentum, in their paper). The conceptual problem is better understood considering the 1951 variant by David Bohm using spin correlations. In the EPR-Bohm experiment the two-particle singlet state can be written as:

$$|\psi\rangle = \frac{1}{\sqrt{2}}[|\uparrow\rangle_1 \otimes |\downarrow\rangle_2 - |\downarrow\rangle_1 \otimes |\uparrow\rangle_2]$$

where $|\uparrow\rangle_j$ ($|\downarrow\rangle_j$) describes the spin state of $j^{th}$ particle ($j = 1, 2$) with spin up (down) respectively. Measurement of the spin on one particle, undetermined prior to the measurement, will “collapse” the wave function to one of the eigenstates and therefore predicts with certainty the outcome of the spin measurement on the second particle without actually doing any measurement. The
important point is that the spin of the second particle in a given direction 
is defined by the choice of the polarizer orientation on the first particle. The 
orientation can be chosen at the "last moment", just prior to the arrival of the 
particle, and cannot be communicated to the second particle system unless 
superluminal signals are invoked. We should conclude that in a way or 
other the second particle carries the information needed to behave correctly 
for any possible choices of the measurement in the system of the first particle. 
Indeed, following EPR, one can define "elements of reality" for spin in $S_x$ and $S_y$ direction for the second particle, 
determined from the spin measurements done on the first particle. But according to QM the 
observables $S_x$ and $S_y$ do not commute and therefore cannot have definite values at the same 
time. EPR-Bohm then concludes that the description of reality given by QM is incomplete. This 
points to the need of extra information, "hidden variables" (HV) for instance, to complement 
QM. In 1964 J. S. Bell found a general scheme to test QM against HV theories: he showed that 
a certain inequality which is always satisfied by all local hidden variable models, can instead be 
violated by QM. Several experiments have been performed, mostly trying to apply a Bell test 
for measurement of the polarization of low energy photons. In the domain of high energy physics, 
CPEar and KLOE have studied correlations in $K^0-\bar{K}^0$ pairs, and obtained results in 
agreement with QM predictions. In this paper we present a study of EPR correlation in the 
flavour of neutral B-meson pairs from $\Upsilon(4S)$ decays. The system is described by a 
wavefunction analogous to (1) 

$$|\psi\rangle = \frac{1}{\sqrt{2}} ([B^0]_1 \otimes \overline{[B^0]}_2 - \overline{[B^0]}_1 \otimes [B^0]_2).$$

Decays occurring at the same proper time are fully correlated: the flavour-specific decay of one 
meson fixes the (previously undetermined) flavour ($B^0$ or $\overline{B^0}$) of the other meson. From (2) we 
deduce the time-dependent rate for decay into two flavour-specific states for opposite flavour (OF, 
$B^0\overline{B^0}$) and same flavour (SF, $B^0B^0$ or $\overline{B^0}\overline{B^0}$) decays, and the corresponding time-dependent 

$$R_{OF} = e^{-\Delta t/\tau_{B^0}}/\langle 4\tau_{B^0} \rangle \{1 \pm \cos(\Delta m_d \Delta t)\},$$

$$A_{QM}(\Delta t) = \frac{R_{OF} - R_{SF}}{R_{OF} + R_{SF}} = \cos(\Delta m_d \Delta t)$$

$\Delta t \equiv |t_1 - t_2|$ is the proper-time difference of the decays, and $\Delta m_d$ the mass difference between 
the two $B^0-\overline{B^0}$ mass eigenstates. We have assumed a lifetime difference $\Delta t_{B^0} = 0$ and neglected the 
$O(10^{-4})$ effects of CP violation in mixing. The fact that the asymmetry depends only on 
$\Delta t$, and not on the absolute time, $t_1$ and/or $t_2$, is a manifestation of EPR-type entanglement at 
a distance. It must be noticed that experimentally it is very difficult to measure the absolute 
times $t_1$ and $t_2$, hence only $\Delta t$ is available.

To be able to reject HV models, ideally a Bell test should be performed. An early attempt in 
this direction was found incorrect. In general Bell tests are unaccessible due to the 
rapid decrease in time of the $B$-meson amplitudes, and the passive character of the flavour 
measurement. Ultima ratio, to probe the non-local behaviour of the $B^0$ pair we can pragmatically 
limit ourselves to verify that, first, QM reproduces the experimental asymmetry, and, second, this 
is not the case for any other "reasonable" HV-based model. Within the definition of "reasonable" 
we include the capability to reproduce the $B^0-\overline{B^0}$ oscillation behaviour for each boson taken 
individually, after the $\Upsilon(4S)$ decay. In conclusion, we have chosen to compare our results with 
the predictions of QM and two other models. We stress the fact that to keep open the possibility 
of testing more models we also provide a fully corrected experimental time-dependent asymmetry, i.e. 
the background is subtracted and the detector effects corrected by deconvolution.

In the local realistic model by Pompili and Selleri (PS), each $B$ transports flavour information ($B^0$ or $\overline{B^0}$), and mass (corresponding to the heavy and light $B_H$, $B_L$ eigenstates). There are
thus four basic states: $B^0_L$, $B^0_{L'}$, $\bar{B}^0_L$, $\bar{B}^0_{L'}$. The model imposes mass and flavour anti-correlations at equal times $\Delta t = 0$; mass values are stable, but the system is programmed to allow random simultaneous jumps in flavour within the pair. The model is also required to reproduce the QM predictions for uncorrelated $B$-decays. No other assumptions are made: the result is an upper and a lower bound for the asymmetry,

$$A_{PS}^{\text{MAX}}(t_1, t_2) = 1 - \{1 - \cos(\Delta m_d \Delta t)\} \cos(\Delta m_d t_{\text{min}}) + \sin(\Delta m_d \Delta t) \sin(\Delta m_d t_{\text{min}}),$$  

$$A_{PS}^{\text{MIN}}(t_1, t_2) = 1 - \min(2 + \Psi, 2 - \Psi), \quad \text{where}$$  

$$\Psi = \{1 + \cos(\Delta m_d \Delta t)\} \cos(\Delta m_d t_{\text{min}}) - \sin(\Delta m_d \Delta t) \sin(\Delta m_d t_{\text{min}}).$$

Note the additional $t_{\text{min}} = \min(t_1, t_2)$ dependence, which can be removed by integrating the OF and SF functions for fixed values of $\Delta t$. We obtain the curves $PS_{\text{max}}$ and $PS_{\text{min}}$ shown in Fig. 1.

In the Spontaneous and immediate Disentanglement model (SD), the $B$-meson pair separates into a $B^0$ and $\bar{B}^0$ with well-defined flavour immediately after the $\Upsilon(4S)$ decay, which then evolve independently, and the asymmetry becomes

$$A_{SD}(t_1, t_2) = \cos(\Delta m_d t_1) \cos(\Delta m_d t_2) = \frac{1}{2}[\cos(\Delta m_d(t_1 + t_2)) + \cos(\Delta m_d \Delta t)],$$

depending on $t_1 + t_2$ in addition to $\Delta t$. After integration we obtain the curve SD of Fig. 1.

Finally, assuming QM as the correct model, we can consider hypothetical effects which can disturb the propagation of the entangled wave function, and affect the time-dependent asymmetry. Suitable parameterisations of the asymmetry for disentanglement in the flavour and mass bases are

$$A = (1 - \zeta_{B^0\bar{B}^0})A_{QM} + \zeta_{B^0\bar{B}^0}A_{SD}, \quad \text{and}$$  

$$A = (1 - \zeta_{B^0L})A_{QM}$$

respectively. In a simplified approach which assumes immediate partial disentanglement into flavour or mass eigenstates, the $\zeta$ parameters correspond to the fraction of decoherent $B$-pairs. (Eq. (10) corresponds to formula 3.5 in Ref. 5, for $\Delta t = 0$).

2 Data analysis

To determine the asymmetry we use $152 \times 10^6 B\bar{B}$ pairs collected by the Belle detector at the $\Upsilon(4S)$ resonance at the KEKB asymmetric-energy (3.5 GeV on 8.0 GeV) $e^+e^-$ collider, by
the Belle detector\(^{16}\). The \(\Upsilon(4S)\) is produced with \(\beta\gamma = 0.425\) close to the \(z\) axis. As the \(B\) momentum is low in the \(\Upsilon(4S)\) center-of-mass system (CMS), \(\Delta t\) can be determined from the \(z\)-displacement of \(B\)-decay vertices: \(\Delta t \approx \Delta z / \beta\gamma c\). The Belle vertex detector provides \(\Delta z\) with a precision of about 100 \(\mu m\).

The event selection for this study (see Ref.\(^{17}\) for details) was optimized for theoretical model discrimination. The flavour of one neutral \(B\) was obtained by reconstructing the decay \(B^0 \rightarrow D^*\ell^+\nu\), with \(D^{*-} \rightarrow D^0\pi^-\) and \(D^0 \rightarrow K^+\pi^-\pi^0\) or \(K^+\pi^-\pi^+\pi^-\) (charge-conjugate modes are included throughout this paper). The \(D^0\) candidates must have a reconstructed mass compatible with the known value. A \(D^*\) is formed by constraining a \(D^0\) and a slow pion to a common vertex. We require a mass difference \(M_{\text{diff}} = M_{K\pi\pi} - M_{K\pi\pi} \in [144.4, 146.4]\) MeV/c\(^2\) (Fig. 2, left), and CMS momentum \(p_{D^*}^+ < 2.6\) GeV/c, consistent with \(B\)-decay. We require that the CMS angle between the \(D^*\) and lepton be greater than 90°. From the relation \(M^2_D = (E^*_B - E^*_D, \ell)^2 - |p^*_B|^2 - |p^*_D, \ell|^2 + 2|p^*_B||p^*_D, \ell|\cos(\theta_{B, D^*}, \ell),\) where \(\theta_{B, D^*}, \ell\) is the angle between \(p^*_B\) and \(p^*_D, \ell\), we can reconstruct \(\cos(\theta_{B, D^*}, \ell)\) by assuming a vanishing neutrino mass. We require \(|\cos(\theta_{B, D^*}, \ell)| < 1.1\). The neutral \(B\) decay position is determined by fitting the lepton track and \(D^0\) trajectory to a vertex, constrained to lie in the \(e^+e^-\) interaction region. The remaining tracks are used to determine the second \(B\) decay vertex and flavour\(^{18}\).

In total 8565 events are selected (6718 OF, 1847 SF). To compensate for the rapid fall in event rate with \(\Delta t\), the time-dependent distributions are histogrammed in 11 variable-size bins (see Table 1). The raw asymmetry is shown in Fig. 2, right. Background subtraction is then performed bin-by-bin; systematic errors are likewise determined by estimating variations in the OF and SF distributions, and calculating the effect on the asymmetry.

A GEANT-based Monte Carlo (MC) sample was analysed with identical criteria, and used for consistency checks, background estimates and subtraction, and to build deconvolution matrices.

Four types of background events have been considered: \(e^+e^- \rightarrow q\bar{q}\) continuum, fake \(D^*\), wrong \(D^*\)-lepton combinations, and \(B^+ \rightarrow D^{*0}\pi^+\pi^-\) events. Off-resonance data (8.3 fb\(^{-1}\)) were used to estimate the continuum background, which was found to be negligible. Fake \(D^0\) reconstruction and misassigned slow pions producing a fake \(D^*\) background were estimated from the sideband in \(M_{\text{diff}}\) (Fig. 2, left). The contamination from wrong \(D^*\)-lepton combinations was obtained by a reverse lepton momentum method, the validity of which was confirmed by MC studies. A fit of the \(\cos(\theta_{B, D^*}, \ell)\) distribution allows the extraction of the \(D^{*+}\) component. The MC is then used to compute the fraction from charged \(B\) mesons which must be subtracted (as it has no mixing).

After correction for wrong flavour assignments (an event fraction of 0.015 ± 0.005) using OF and SF distributions from wrongly-tagged MC events, we obtain the time-dependent asymmetry.
Table 1: Time-dependent asymmetry in $\Delta t$ bins, corrected for experimental effects, with total uncertainties.

<table>
<thead>
<tr>
<th>bin</th>
<th>window $[\text{ps}]$</th>
<th>$A$ and total error</th>
<th>bin</th>
<th>window $[\text{ps}]$</th>
<th>$A$ and total error</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>$0.0 - 0.5$</td>
<td>$1.013 \pm 0.028$</td>
<td>7</td>
<td>$5.0 - 6.0$</td>
<td>$-0.961 \pm 0.077$</td>
</tr>
<tr>
<td>2</td>
<td>$0.5 - 1.0$</td>
<td>$0.916 \pm 0.022$</td>
<td>8</td>
<td>$6.0 - 7.0$</td>
<td>$-0.974 \pm 0.080$</td>
</tr>
<tr>
<td>3</td>
<td>$1.0 - 2.0$</td>
<td>$0.699 \pm 0.038$</td>
<td>9</td>
<td>$7.0 - 9.0$</td>
<td>$-0.675 \pm 0.109$</td>
</tr>
<tr>
<td>4</td>
<td>$2.0 - 3.0$</td>
<td>$0.339 \pm 0.056$</td>
<td>10</td>
<td>$9.0 - 13.0$</td>
<td>$0.089 \pm 0.193$</td>
</tr>
<tr>
<td>5</td>
<td>$3.0 - 4.0$</td>
<td>$-0.136 \pm 0.075$</td>
<td>11</td>
<td>$13.0 - 20.0$</td>
<td>$0.243 \pm 0.435$</td>
</tr>
<tr>
<td>6</td>
<td>$4.0 - 5.0$</td>
<td>$-0.634 \pm 0.084$</td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

shown in Fig. 2, right.

Remaining experimental effects (e.g. resolution in $\Delta t$, selection efficiency) are corrected by a deconvolution procedure\textsuperscript{19}. $11 \times 11$ response matrices are built separately for SF and OF events, using MC $D^*\ell\nu$ events indexed by generated and reconstructed $\Delta t$ values. The procedure has been optimised, and its associated systematic errors inferred by a toy Monte Carlo where sets of several hundred simulated experiments are generated assuming the three theoretical models. We test the consistency of the method applied to our data by fitting the $B^0$ decay time distribution (summing OF and SF samples), leaving the $B^0$ lifetime as a free parameter. We obtain $1.532 \pm 0.017$ (stat) ps, consistent with the world average\textsuperscript{20}. We have also repeated the deconvolution procedure using a subset of events with better vertex fit quality, and hence more precise $\Delta t$ values: consistent results are obtained. The final results are shown in Table 1 and Fig. 3.

Figure 3: Bottom: time-dependent flavour asymmetry (crosses) and the results of weighted least-squares fits to the (left to right) QM, SD, and PS models (rectangles, showing $\pm 1\sigma$ errors on $\Delta m_d$). Top: differences $\Delta \equiv A_{\text{data}} - A_{\text{model}}$ in each bin, divided by the total experimental error $\sigma_{\text{tot}}$. Bins where $A_{\text{PS}}^{\text{min}} < A_{\text{data}} < A_{\text{PS}}^{\text{max}}$ have been assigned a null deviation: see the text.

3 Comparison with the theoretical models

The model testing is done by a least-square fit to $A(\Delta t)$, leaving $\Delta m_d$ free, but taking the world-average $\Delta m_d$ into account. To avoid bias, we discard BaBar and Belle measurements, which assume QM correlations: this yields\textsuperscript{21} $(\Delta m_d) = (0.496 \pm 0.014)$ ps$^{-1}$. Our data is in agreement with the prediction of QM: we obtain $\Delta m_d = 0.501 \pm 0.009$ ps$^{-1}$ with $\chi^2 = 5.2$ for 11 dof (see Fig. 3). SD is rejected by $\chi^2 = 174$ ($\Delta m_d = 0.419 \pm 0.008$). To fit PS we have used the closest boundary to our data $A_{\text{PS}}^{\text{max}}$, Eq. (5), or $A_{\text{PS}}^{\text{min}}$, Eq. (6), but assumed a null deviation
for data falling inside the boundaries. We obtain $\chi^2 = 31.3$ ($\Delta m_d = 0.447 \pm 0.010 \text{ ps}^{-1}$): the data favour QM over PS at the 5.1$\sigma$ level.

We have examined the possibility of a partial loss of coherence just after the decay of the $\Upsilon(4S)$ resonance. The fraction of events with disentangled $B^0$ and $\overline{B^0}$ can be estimated by fitting our asymmetry with the mixture of Eq. (9), leaving $\zeta_{B^0\overline{B^0}}$ free. The fit finds $\zeta_{B^0\overline{B^0}} = 0.029 \pm 0.057$, consistent with no decoherence. The second possibility considered is a decoherence into mass eigenstate, for which we expect a reduction in the amplitude of $A(\Delta t)$, as given by Eq. (10). The result of a fit gives a value of $\zeta_{B^0\overline{B^0}_L} = 0.004 \pm 0.017$ (preliminary), also compatible with zero.

4 Conclusion

We have analysed neutral $B$ pairs produced by $\Upsilon(4S)$ decay, determined the time-dependent asymmetry due to flavour oscillations, and corrected for experimental effects by deconvolution: the results can be directly compared to theoretical models. Given the fact that there is little hope to perform a Bell test in the neutral $B$ system, we have compared our data to the QM hypothesis and to two other models. The local realistic model of Pompili and Selleri is strongly disfavoured compared to the entanglement predicted by QM. Immediate disentanglement, in which definite-flavour $B^0$ and $\overline{B^0}$ evolve independently, is ruled out. We have also found that our data is consistent with a null fraction of events with a loss of entanglement.

References

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OBSERVATION OF THE Σ₀ BARYONS AT CDF

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We present a measurement of four new bottom baryons in proton-antiproton collisions with a center of mass energy of 1.96 TeV. Using 1.1 fb⁻¹ of data collected by the CDF II detector, we observe four Λ₀ b→Λ⁺ b mark resonances in the fully reconstructed decay mode Λ₀ b→Λ⁺ b π⁻, where Λ⁺ b→pK⁺π⁻. The probability for the background to produce a similar or larger signal is less than 8.3 x 10⁻⁶, corresponding to a significance of greater than 5.2 σ. We interpret these baryons as the Σ₀ b mark states.

1 Introduction

The Tevatron at the Fermi National Accelerator Laboratory collides p̅p with a center of mass energy of 1.96 TeV. The Collider Detector at Fermilab, or CDF, experiment employs a general multipurpose detector to reconstruct particle physics events from these collisions. With a b hadron cross section of ≃50 µb (|η| < 1.0), CDF has collected a wealth of experimental data on b hadrons. Using this data, we announce the first observation of the Σ₀ b mark baryons.

2 Σ₀ b mark Theoretical Predictions

Only one b baryon has been previously established, the ground state Λ₀ b, which contains b, u, and d quarks with the two light quarks (u and d) in a flavor antisymmetric diquark state. CDF uses a two displaced track trigger to select the decay of Λ₀ b→Λ⁺ bπ⁻, with Λ⁺ b→pK⁺π⁻ (inclusion of the respective charge conjugate modes is assumed throughout this paper). The two displaced track trigger requires two high p_T tracks displaced from the p̅p interaction point; in the decay of Λ₀ b, the two tracks which satisfy the requirements are primarily the pion from the Λ₀ b decay and the proton from the Λ⁺ b decay. Using 1.1 fb⁻¹ of data collected by the CDF II detector between February 2002 and March 2006, CDF possesses the world’s largest sample of bottom baryons with 3180 ± 60 (stat.) Λ₀ b→Λ⁺ bπ⁻ candidates. The reconstructed Λ₀ b invariant mass distribution is shown in Fig. 1.
Figure 1: Unbinned maximum likelihood fit to the reconstructed invariant mass of $\Lambda_{\mu}^{0} \rightarrow \Lambda_{\mu}^{0} \pi^{-}$ candidates. The fully reconstructed $\Lambda_{\mu}^{0}$ modes (such as $\Lambda_{\mu}^{0} \rightarrow \Lambda_{\mu}^{0} \pi^{-}$ and $\Lambda_{\mu}^{0} \rightarrow \Lambda_{\mu}^{0} K^{-}$) are not shown separately on the figure.

Figure 2: Unbinned maximum likelihood fit to the $\Sigma_0$ $Q$-distributions. The top plot shows the $\Lambda_{\mu}^{0} \Xi_{c}^{0}$ combinations, while the bottom plot shows the $\Lambda_{\mu}^{0} \Xi^{-}$ combinations. The insets show the expected background plotted on the data for $Q \in [0, 500]$ MeV/$c^2$, while the $\Sigma_0$ signal fit is shown on a reduced range of $Q \in [0, 200]$ MeV/$c^2$. 
The next accessible \( b \) baryons are the lowest lying \( \Sigma_b^{||} \) states, which decay strongly to \( \Lambda_b^0 \) baryons by emitting pions. The \( \Sigma_b^{+} \) baryons contain one \( b \) and two \( u \) quarks, the \( \Sigma_b^- \) baryons contain one \( b \) and two \( d \) quarks, and the \( \Sigma_b^{0} \) baryons contain the \( b, u, \) and \( d \) quarks. In the \( \Sigma_b \) baryons, the two light quarks are in a flavor symmetric diquark state, leading to a doublet of baryons with \( J^P = \frac{1}{2}^+ \) (\( \Sigma_b \)) and \( J^P = \frac{3}{2}^+ \) (\( \Sigma_b^* \)). Because \( \Sigma_b^{0,\pm} \) decays to \( \Lambda_b^0 \pi^0 \) and the CDF II detector cannot reconstruct neutral pions, we expect to observe only the charged \( \Sigma_b^{0,\pm} \) states.

There is predicted to be a hyperfine mass splitting between the doublet states \( \Sigma_b \) and \( \Sigma_b^* \), as well as a mass splitting between the \( \Sigma_b^{||} \) and \( \Sigma_b^{+} \) states due to strong isospin violation. Predictions for the \( \Sigma_b^{||} \) masses exist from heavy quark effective theories, non-relativistic and relativistic potential models, \( 1/N_c \) expansion, sum rules, and lattice Quantum Chromodynamics calculations. These predictions expect \( m(\Sigma_b) - m(\Lambda_b^0) \approx 180 - 210 \text{ MeV/c}^2 \), \( m(\Sigma_b^0) - m(\Sigma_b) \approx 10 - 40 \text{ MeV/c}^2 \), and \( m(\Sigma_b^-) - m(\Sigma_b^0) \approx 5 - 7 \text{ MeV/c}^2 \). The intrinsic width of \( \Sigma_b^{0,\pm} \) baryons is dominated by the P-wave one-pion transition, whose partial width depends on the available phase space. For the predicted range of \( \Sigma_b^{||} \) masses, the intrinsic width varies between 2 and 20 MeV/c^2.

3 Analysis Methodology

We search for four resonant \( \Lambda_b^0 \pi^\pm \) states consistent with theoretical predictions for \( \Sigma_b \), where \( \Sigma_b \) now refers to both charged \( J^P = \frac{1}{2}^+ \) (\( \Sigma_b^0 \)) and \( J^P = \frac{3}{2}^+ \) (\( \Sigma_b^{+} \)) states. To minimize the contribution of the mass resolution of each \( \Lambda_b^0 \) candidate, the search is made for narrow resonances in the mass difference distribution of \( Q = m(\Lambda_b^0 \pi^-) - m(\Lambda_b^0) - m(\pi^-) \). Events are separated into "\( \Lambda_b^0 \pi^- \)" and "\( \Lambda_b^0 \pi^+ \)" subsamples; \( \Lambda_b^0 \pi^- \) contains \( \Sigma_b^{||} \) and \( \Sigma_b^- \) while \( \Lambda_b^0 \pi^+ \) contains \( \Sigma_b^{+} \) and \( \Sigma_b^{||} \).

The \( \Lambda_b^0 \) candidate is combined with a prompt pion, as the \( \Sigma_b \) decays strongly at the primary vertex of the \( p\bar{p} \) collision. To perform an unbiased optimization of the selection criteria, we use as a background sample only those tracks far from the expected \( \Sigma_b \) signal region. From theoretical predictions, the signal region is defined as \( 30 < Q < 100 \text{ MeV/c}^2 \). The principle sources of background in the \( \Sigma_b \) \( Q \) distribution are tracks from the hadronization of prompt \( \Lambda_b^0 \) baryons and \( B \) mesons reconstructed as \( \Lambda_b^0 \) baryons, and combinatorial background. The percentage of each background source in the \( \Sigma_b \) \( Q \) distribution is fixed from the \( \Lambda_b^0 \) invariant mass fit shown in Fig. 1, which is 89.4% \( \Lambda_b^0 \) baryons, 7.3% \( B \) mesons, and 3.3% combinatorial background. The \( Q \) distribution of each background component is established before unblinding the signal region. The high mass region above the \( \Lambda_b^0 \to \Lambda_b^+ \pi^- \) signal in the \( \Lambda_b^0 \) mass distribution (Fig. 1) determines the combinatorial background. Reconstructing \( B^0 \to D^+ \pi^- \) data as \( \Lambda_b^0 \to \Lambda_b^+ \pi^- \) gives the background from \( B \) hadronization tracks. The largest background component, from \( \Lambda_b^0 \) hadronization tracks, is obtained from a \( \Lambda_b^0 \) PYTHIA Monte Carlo simulation.

4 \( \Sigma_b^{||} \) Results

After determining the background shape, we observe an excess of events over the expected background in the \( \Sigma_b \) signal region. The excess in the \( \Lambda_b^0 \pi^- \) subsample is 118 candidates over 288 expected background candidates, while in the \( \Lambda_b^0 \pi^+ \) subsample the excess is 91 over 313 expected background candidates. The strength of the \( \Sigma_b \) hypothesis is evaluated using as a test statistic the likelihood ratio, \( LR = L_{\mathrm{alt}} / L \), where \( L \) is the fit likelihood of the four \( \Sigma_b \) signal hypothesis and \( L_{\mathrm{alt}} \) is the fit likelihood of an alternate hypothesis such as the no signal hypothesis. Using simplistic Monte Carlo samples of background fluctuations, we find the probability of the no signal hypothesis to be less than \( 8.3 \times 10^{-8} \), corresponding to a signal significance of greater than 5.2 \( \sigma \).

The subsamples are modeled with a simultaneous unbinned maximum likelihood fit comprising a signal for each expected \( \Sigma_b \) state plus the background. Each signal consists of a non-relativistic Breit-Wigner distribution convoluted with a double Gaussian model of the detector resolution. The intrinsic
width of the Breit-Wigner is computed from the available phase space given the central location of the signal. Due to low statistics, the constraint \( m(\Sigma_b^{+}) - m(\Sigma_b^{-}) = m(\Sigma_b^{0}) - m(\Sigma_b^{+}) = \Delta(\Sigma_b^{0}) \) is added. The \( \Sigma_b \) signal fit to data, which has a \( \chi^2 \) fit probability of 76% in the range \( Q \in [0, 200] \) MeV/c², is shown in Fig. 2. The majority of the systematic uncertainty on the yield measurement is due to poor knowledge of the \( \Lambda_b \) hadronization background, while the majority of the systematic uncertainty on the mass measurement is due to the CDF mass scale uncertainty. The final results for the yields are \( N(\Sigma_b^{+}) = 32^{+13}_{-12} \) (stat.)\( ^{+5}_{-3} \) (syst.), \( N(\Sigma_b^{0}) = 59^{+15}_{-14} \) (stat.)\( ^{+9}_{-2} \) (syst.), \( N(\Sigma_b^{+}) = 77^{+17}_{-16} \) (stat.)\( ^{+10}_{-7} \) (syst.), and \( N(\Sigma_b^{-}) = 69^{+18}_{-17} \) (stat.)\( ^{+16}_{-5} \) (syst.). The signal locations are \( Q(\Sigma_b^{+}) = 48.5^{+2.4}_{-2.2} \) (stat.)\( ^{+0.2}_{-0.3} \) (syst.) MeV/c², \( Q(\Sigma_b^{0}) = 55.9 \pm 1.0 \) (stat.)\( \pm 0.2 \) (syst.) MeV/c², and \( \Delta(\Sigma_b^{0}) = 21.2^{+2.0}_{-1.9} \) (stat.)\( ^{+0.4}_{-0.3} \) (syst.) MeV/c².

5 Summary

The \( \Lambda_b^0 \pi^\pm \) resonant states observed in 1.1 fb\(^{-1}\) of CDF II data are consistent with the lowest lying charged \( \Sigma_b \) baryons, and the observed properties are in agreement with theoretical predictions. Using the CDF II measurement\(^6\) of \( m(\Lambda_b^0) = 5619.7 \pm 1.2 \) (stat.)\( \pm 1.2 \) (syst.) MeV/c², the masses of each state are

\[
\begin{align*}
    m(\Sigma_b^{+}) &= 5807.8^{+2.0}_{-2.2} \, \text{MeV/c}^2, \\
    m(\Sigma_b^{0}) &= 5815.2 \pm 1.0 \pm 1.7 \, \text{MeV/c}^2, \\
    m(\Sigma_b^{+}) &= 5829.0^{+1.6}_{-1.8} \, \text{MeV/c}^2, \\
    m(\Sigma_b^{0}) &= 5836.4 \pm 2.0 \pm 1.8 \, \text{MeV/c}^2.
\end{align*}
\]

This is the first observation of the lowest lying charged \( \Sigma_b \) baryons.

Acknowledgments

We thank T. Becher, A. Falk, D. Pirjol, J. Rosner, and D. Ebert for useful discussions. We also thank the Fermilab staff and the technical staffs of the participating institutions for their vital contributions. This work was supported by the U.S. Department of Energy and National Science Foundation; the Italian Istituto Nazionale di Fisica Nucleare; the Ministry of Education, Culture, Sports, Science and Technology of Japan; the Natural Sciences and Engineering Research Council of Canada; the National Science Council of the Republic of China; the Swiss National Science Foundation; the A.P. Sloan Foundation; the Bundesministerium für Bildung und Forschung, Germany; the Korean Science and Engineering Foundation and the Korean Research Foundation; the Particle Physics and Astronomy Research Council and the Royal Society, UK; the Institut National de Physique Nucleaire et Physique des Particules/CNRS; the Russian Foundation for Basic Research; the Comisión Intermínisterial de Ciencia y Tecnologia, Spain; the European Community’s Human Potential Programme under contract HPRN-CT-2002-00292; and the Academy of Finland.

References

2. The CDF II detector uses a cylindrical coordinate system with the z-axis along the proton beam direction. The pseudorapidity \( \eta \) is defined as \( \tanh^{-1}(\cos \theta) \). Transverse momentum, \( p_T \), is the component of the particle’s momentum in the \( \{x,y\} \) plane.
Searching for $H \to WW^*$ and Other Di-boson Final States at CDF

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La Jolla, California 92093

We report a search for standard model (SM) production of a Higgs boson which decays to $WW^*$ in two charged leptons ($e, \mu$) and two neutrino final state in $p\bar{p}$ collisions at $\sqrt{s} = 1.96$ TeV. The data were collected with the CDF II detector at the Fermilab Tevatron and correspond to an integrated luminosity of 1.1 fb$^{-1}$. The Matrix Element method is used to calculate the event probability and to construct a likelihood ratio discriminator. The observed (the median of the expected) 95% Confidence Level (CL) upper limit for $\sigma(H \to WW^*)$ with 100 GeV/c$^2$ mass hypothesis is 1.3(1.8) pb which is 3.4(4.8) times the SM prediction at next-to-next-to-leading logarithmic level (NNLL) calculation. The SM ZZ production search is performed in the same final states. The observed significance is 1.9 σ and the 95% CL upper limit is 3.4 pb which is consistent with the next-to-leading order(NLO) calculation of 1.4 ± 0.1 pb.

1 Introduction

The Higgs boson is introduced into the standard model (SM) to explain the electroweak symmetry breaking and the origins of particle mass. The 95% CL interval on the Higgs mass is constrained to be 114-182 GeV/c$^2$ with the current precision electroweak measurements. We search for the Higgs through gluon fusion production channel, $gg \to H \to WW^*$, which is the dominant channel for Higgs $m_H > 135$ GeV/c$^2$. The maximum NNLL Higgs production cross section is $\sigma(pp \to H \to WW^*) = 0.388$ pb for $m_H = 160$ GeV/c$^2$. This is a small signal compared to the SM $WW$ production with NLO cross section 12.4 pb. A good understanding of the SM diboson production is essential for this search. To get a good signal to background ratio sample, we search for fully leptonic decay of $WW^* \to l^+l^-\nu\bar{\nu}$, where $l^\pm = e, \mu$ or $\tau$ and $\tau$ decays to $e$ or $\mu$. The ZZ could decay to the same final states and it has not yet been observed at $p\bar{p}$ colliders. The analysis strategy is to maximize the signal acceptance by loosing selection cuts and use the likelihood ratio discriminator (LR) calculated by Matrix Element methods to
set the statistical limits for 10 different Higgs mass hypothesis. The search for ZZ production is done in the same way but considering ZZ as signal and no Higgs contribution.

2 Selection

The $e^{+}e^{-}\nu\bar{\nu}$ candidates are selected from two opposite-sign leptons from the same vertex and high missing transverse energy $H_{T}$. At least one lepton is required to satisfy the trigger and have $p_T > 20$ GeV/c. The other lepton has looser requirement $p_T > 10$ GeV/c to increase the kinematic acceptance. To suppress the significant backgrounds from $W\gamma$ and $W \pm$ jets where $\gamma$ conversions to electrons or a jet is mis-constructed as a lepton, we require leptons to be both energy and track isolated such that the sum of the $E_T(p_T)$ for the calorimeter towers (tracks) in a cone of $\Delta R = \sqrt{(\Delta \eta)^2 + (\Delta \phi)^2} < 0.4$ around the lepton is less than 10% of the $E_T$ for electrons or $p_T$ for muons and track lepton candidates. To suppress the Drell-Yan background, we require $\min H_{T,rel} > 25$ GeV, where $\min H_{T,rel}$ is defined to be:

$$\min H_{T,rel} \equiv \begin{cases} H_{T} & \text{if } \Delta \phi(H_{T}, \text{lepton, jet}) > \frac{\pi}{2} \\ H_{T} \sin(\Delta \phi(H_{T}, \text{lepton, jet})) & \text{if } \Delta \phi(H_{T}, \text{lepton, jet}) < \frac{\pi}{2} \end{cases}$$

(1)

This definition will reject the events whose $H_{T}$ just comes from single lepton or jet. We further require the candidates to have less than 2 jets with $p_T > 15$ GeV and $|\eta| < 2.5$, in order to suppress $W$ backgrounds, $M_{T1,T2} > 25$ GeV in order to suppress heavy flavor contributions, and exactly 2 leptons to suppress $WZ$ contributions with a third lepton.

For ZZ analysis, the $e\mu$ channel is not used and one addition cut, $H_{T,\text{sig}} \equiv H_{T} \sqrt{\sum E_T} > 2.5$ GeV, is applied to suppress the effect of mis-measurement of unclustered energy.

3 Event Proximity Calculation

In order to use the maximum kinematic information to distinguish each mode, we use an event-by-event calculation of the probability density function $P_m(x_{obs})$ for a mode $m$ which is either Higgs, WW, ZZ, W$\gamma$ or $W$+parton:

$$P_m(x_{obs}) = \frac{1}{\sigma_m} \int \frac{d \sigma_m^{th}(y)}{dy} \epsilon(y) G(x_{obs}, y) dy$$

(2)

where $x_{obs}$ are the observed leptons four-vectors and $H_{T}$, $y$ are the true lepton four-vectors (include neutrinos), $\sigma_m^{th}$ is the MCFAM leading-order theoretical calculation of the cross-section for mode $m$, $\epsilon(y)$ is total event efficiency $\times$ acceptance, $G(x_{obs}, y)$ is an analytic model of resolution effects, and $\frac{1}{\sigma_m}$ is the normalization. The function $\epsilon(y)$ describes the probabilities of a parton level object ($e$, $\mu$, $\gamma$ or parton) to be reconstructed as an observed lepton and is extracted by combinations of Monte Carlo and data. The event probability density functions are used to construct one dimensional discriminator:

$$LR(x_{obs}) = \frac{P_H(x_{obs})}{P_H(x_{obs}) + \Sigma_k k_k P_k(x_{obs})},$$

(3)

where $H$ is Higgs, $k_k$ is the expected fraction for each background and $\Sigma_k k_k = 1$. For SM ZZ search, we just use ZZ and WW to construct the discriminator.

4 Systematics

Table 1 summarizes the various systematics of each mode. The $H_{T}$ resolution modeling uncertainty, lepton selection scale factor and trigger efficiency are determined from comparisons of
the data and the Monte Carlo simulation in a sample of dilepton events. The uncertainties are propagated through the analysis. For the $W\gamma$ background contribution, there is an additional uncertainty of 20% from the detector material description and conversion veto efficiency. The higher order effects in $WW$ is assigned to be a half of the difference between the Pythia and MC@NLO acceptance. The systematic uncertainty on the $W$+jets background is determined from differences between the measured probability that a jet is identified as a lepton for jets collected using different jet $E_T$ trigger thresholds. An additional 6% uncertainty originating from the luminosity measurement is assigned to both signal and background except $W$+jets.

<table>
<thead>
<tr>
<th>$E_T$ Modeling</th>
<th>$WW$</th>
<th>$WZ$</th>
<th>$ZZ$</th>
<th>$tt\bar{t}$</th>
<th>$W\gamma$</th>
<th>$W$+jets</th>
<th>Higgs</th>
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</thead>
<tbody>
<tr>
<td>NLO Acceptance</td>
<td>4.5 (5.1)</td>
<td>10.0</td>
<td>10.0</td>
<td>10.0</td>
<td>5.0</td>
<td>10.0</td>
<td>10.0</td>
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<td>Cross-section</td>
<td>10.0</td>
<td>10.0</td>
<td>10.0</td>
<td>15.0</td>
<td>5.0</td>
<td>10.0</td>
<td>10.0</td>
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<tr>
<td>PDF Uncertainty</td>
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<td>2.7</td>
<td>3.1</td>
<td>4.3</td>
<td>2.2</td>
<td>2.2</td>
</tr>
<tr>
<td>$\text{Log}\frac{d^2\sigma}{d^2E_T}$</td>
<td>1.2 (1.4)</td>
<td>1.4 (1.5)</td>
<td>1.4 (1.5)</td>
<td>1.4 (1.5)</td>
<td>1.4 (2.2)</td>
<td>1.4 (1.3)</td>
<td>1.5</td>
</tr>
<tr>
<td>Trigger Eff</td>
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<td>0.3</td>
<td>0.3</td>
<td>0.4</td>
<td>0.4</td>
<td>0.5 (0.6)</td>
<td>0.3</td>
</tr>
<tr>
<td>Total</td>
<td>11.3 (11.5)</td>
<td>14.5</td>
<td>14.5</td>
<td>18.0</td>
<td>21.7</td>
<td>24.7 (24.6)</td>
<td>26.8 (23.3)</td>
</tr>
</tbody>
</table>

5 $H \rightarrow WW^*$ Results

![Figure 1: The LR distributions of Higgs mass 160 GeV/$c^2$ for (a) High S/B channel and (b) Low S/B channel.](image)

The expected and observed yields of each mode are shown in Table 2 and Table 3. The LR distributions are shown in Fig 1 and all the candidates are cataloged into two channels based on the signal to background ratio (S/B) for each event. The limit of Higgs production cross section is evaluated by performing a Bayesian binned maximum likelihood fit. All of the background normalizations are free parameters in the fit but constrained to their expectations with a set of Gaussian constraints considering all of the assumed correlations between the systematics uncertainties. The limits of Higgs production cross section times $WW^*$ decay branching ratio, $\sigma_{95\%}$, and their ratios to NNLL calculations ($\sigma_{SM}$) are shown in Table 3 and Figure 2.

6 ZZ Results

The expected and observed yields after the ZZ selection are shown in Table 4. The variable, $\log_{10}(1 - LR)$, is used to set the upper limit and is shown in Fig 3. The Frequentist approach is used by performing background-only Monte Carlo experiments based on the expected yields varied within the assigned systematics. For each experiment a test statistic is formed from the difference in the log likelihood value with the background-only model and with the signal yield.
at the best fit value. The observed significance is 1.9σ and we set the 95% CL upper limit of 3.4 pb, which is consistent with the SM NLO cross section of 1.4 ± 0.1 pb.

![Expected yields](image1)

![Expected yields](image2)

### Table 4: Expected and observed yields for ZZ selection.

<table>
<thead>
<tr>
<th>Process</th>
<th>Expected</th>
<th>Observed</th>
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</thead>
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<td>WW</td>
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<td>79.2</td>
</tr>
<tr>
<td>WZ</td>
<td>7.1</td>
<td>5.1</td>
</tr>
<tr>
<td>ZZ</td>
<td>10.7</td>
<td>11.2</td>
</tr>
<tr>
<td>tt</td>
<td>5.1</td>
<td>4.1</td>
</tr>
<tr>
<td>DY</td>
<td>24.0</td>
<td>23.2</td>
</tr>
<tr>
<td>Wγ</td>
<td>13.6</td>
<td>12.6</td>
</tr>
<tr>
<td>W+jets</td>
<td>23.2</td>
<td>23.2</td>
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<tr>
<td>Total</td>
<td>152.9±11.6</td>
<td>152.9±11.6</td>
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<tr>
<td>Data</td>
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<td>182</td>
</tr>
</tbody>
</table>

### References


Search for Neutral Higgs Boson Production in the Decay $h \rightarrow \tau\mu\tau$ with the DØ Detector

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A search for the production of neutral Higgs bosons decaying into $\tau^+\tau^-$ final states is presented. One of the two $\tau$ leptons is required to decay into a muon. The data were collected by the DØ detector and correspond to an integrated luminosity of about 1.0 $fb^{-1}$. No excess is observed above the expected backgrounds. The results are interpreted in the Minimal Supersymmetric Standard Model. In the mass range $90 < m_A < 200$ GeV values of $\tan\beta$ larger than 40-60 are excluded for the no-mixing and the $m_h^{max}$ benchmark scenarios.

1 Introduction

The contribution of $\tau^+\tau^-$ final states from the Standard Model (SM) Higgs production is too small to play any role in the SM Higgs searches in $p\bar{p}$ collision at the Tevatron due to the large irreducible background from $Z \rightarrow \tau^+\tau^-$ production. This is different in the Minimal Supersymmetric Standard Model (MSSM), which introduces two Higgs doublets leading to five Higgs bosons: a pair of charged Higgs boson ($H^\pm$); two neutral CP-even Higgs bosons ($h,H$) and a CP-odd Higgs boson ($A$). At tree level, the Higgs sector of the MSSM is fully described by two parameters, which are chosen to be the mass of the CP-odd Higgs, $m_A$, and the ratio of the vacuum expectation values of the two Higgs doublets, $\tan\beta$. The Higgs production cross-section is enhanced in the region of high $\tan\beta$. In the low $m_A$, high $\tan\beta$ region of the parameter space, Tevatron searches can therefore probe several MSSM benchmark scenarios extending the search.
regions covered by LEP. Inclusive searches for $\phi(= H, h, A) \rightarrow \tau\tau$ have been performed with integrated luminosities of $L = 350 \text{ pb}^{-1}$ by DØ and $L = 310 \text{ pb}^{-1}$ by CDF. Both searches require at least one $\tau$ lepton to decay into an electron ($\tau_e$) or a muon ($\tau_\mu$). In this analysis, only the decay $\phi \rightarrow \tau_\mu \tau$ is considered using an integrated luminosity of $L = 1.0 \text{ fb}^{-1}$. CDF has also recently released a preliminary result with $L = 1.0 \text{ fb}^{-1}$. The biggest improvement in sensitivity compared to the previous analysis comes from using a neural network to improve the separation between signal and background.

2 Event Preselection

The preselection requires one isolated muon with $p_T^\mu > 15 \text{ GeV}$. The event is required to have no other muon that is matched to a track in the central detector with $p_T^\mu > 10 \text{ GeV}$.

Hadronically decaying taus are characterized by a narrow isolated jet that is associated with three or less tracks. Three types of hadronically decaying taus are distinguished:

Type 1: Calorimeter energy cluster, with one associated track and no electromagnetic sub-cluster. This corresponds mainly to the decay $\tau^\pm \rightarrow \pi^\pm \nu$.

Type 2: Calorimeter energy cluster, with one associated track and at least one electromagnetic sub-cluster. This corresponds mainly to the decay $\tau^\pm \rightarrow \pi^\pm \pi^0 \nu$.

Type 3: Calorimeter energy cluster, with three associated tracks, with an invariant mass below 1.7 GeV. This corresponds mainly to the decays $\tau^\pm \rightarrow \pi^\pm \pi^\pm \pi^\mp \nu$.

Tau decays into electrons are usually reconstructed as type-2 taus. These are not removed from the sample. The event is required to contain a $\tau$ candidate at a distance $\Delta R > 0.5$ from the muon direction. The charge of the $\tau$ candidate must be opposite to the muon charge. The transverse momentum $p_T^\tau$ of the $\tau$ candidate measured in the calorimeter must be greater than 15 GeV for $\tau$-type 1 and 2, and greater than 20 GeV for $\tau$-type 3. At the same time the transverse momentum of the track associated with the $\tau$ candidate is required to be $p_T > 15 \text{ GeV}$ for $\tau$-type 1 and $p_T > 5 \text{ GeV}$ for $\tau$-type 2. In the case of $\tau$-type 3, the scalar sum of the transverse momenta of all associated tracks must be greater than 15 GeV.

3 $W \rightarrow \mu \nu$ and Multi-jet Background Estimation

The shape of the $W \rightarrow \mu \nu +$ jet background distribution, where the jet is misidentified as a tau, was simulated using PYTHIA. The normalization, however, was obtained using data.

A contribution to the background is expected from heavy flavour multi-jet events, where a muon from a semi-leptonic decay passes the isolation requirement and a jet is mis-identified as a $\tau$ candidate. In addition, a contribution is expected from light quark multi-jet events where the jets fake both the tau and the muon. The multi-jet background shape is taken from events with at least one muon and one $\tau$ candidate where the muon failed the calorimeter isolation requirement. The normalization of this semi-isolated sample was obtained in a multi-jet enriched sample.

4 Final Event Selection

A set of neural networks, one for each tau type, has been developed to separate the tau leptons from jets. These neural networks make use of input variables that exploit the tau signature such as longitudinal and transverse shower shapes and isolation in the calorimeter and the tracker. The neural network is trained using tau MC events as signal and multi-jet events from data as background to produce a variable that peaks near one for real taus and zero for jets. The tau
candidate is required to have a neural network output greater than 0.9. In the case of type-3 taus this is tightened to 0.95 due to the larger multi-jet background.

It is also possible for muons to fake type-1 or type-2 tau candidates. These fakes are removed by ensuring that type-1 or type-2 tau candidates do not match to a reconstructed muon within a cone of radius $\Delta R_{\mu\tau} = 0.5$.

After selecting events with a high neural network output, there is still a considerable amount of background from $W + \text{jets}$ production. To remove these events, the reconstructed $W$ boson mass, $M_W = \sqrt{2E_T^{\mu}E_T^{\nu}(1 - \cos\Delta\phi)}$ is used, where $E_T^{\mu} = E_{T\mu}p_T^\mu/p_T^\mu$ is the estimated neutrino energy, calculated using the ratio of the muon momentum $p_T^\mu$ and muon transverse momentum $p_T^\mu$. For real $W$ boson decays, this variable peaks near the $W$ boson mass, whereas for the signal and the $Z \rightarrow \tau\tau$ background the variable peaks at zero. Events with $M_W > 20$ GeV are rejected.

To achieve the best separation of the signal from background, neural networks were trained for different Higgs mass points using kinematical variables. The distribution of the visible mass $M_{\text{vis}}$ and the optimised neural networks for a Higgs mass of 160 GeV is shown in Figure 1. There is good agreement between the background expectation and the data.

![Figure 1: Distribution of a) the visible mass $M_{\text{vis}}$ and b) neural network output distribution for a Higgs mass of 160 GeV with all selections applied. The data, shown with error bars, are compared to the sum of the expected backgrounds. Overflow events are added to the last bin. Also shown, in light green, is the signal for a Higgs mass of 160 GeV, normalized to a cross-section of 10 pb. The systematic uncertainty on the background normalisation is 10% and is shown by the shaded area.](image)

5 Results and Conclusion

Limits on the cross-section for Higgs boson production times the branching fraction into tau leptons are derived at 95% Confidence Level (CL). The output from the neural networks, shown in Figure 1 for one mass point, are used in the limit calculation. The distributions for the three tau types are used separately. The cross-section limits are calculated with the CL$_S$ method.

There are various sources of systematic uncertainties that affect signal and background. The most important are the uncertainty on the integrated luminosity (6.1%), the trigger efficiency (3%), the tau energy scale (1 – 11%), the uncertainty in the signal acceptance due to choice of parton distribution function (3.9 – 4.6%), the uncertainty of the tau track matching efficiency (4%), the uncertainty on the tau reconstruction efficiency (3%), the theoretical uncertainty on the $Z$ cross-section (5%) and the uncertainty on the modeling of the multi-jet background (3%). All systematic uncertainties are included in the calculation of the expected and observed limits, assuming 100% correlation between signal and background where appropriate. The expected and observed limits are shown in Figure 2 as a function of the hypothetical Higgs mass.

In the MSSM, the masses and couplings of the Higgs bosons depend on $\tan\beta$ and $m_A$ at tree level. Radiative corrections introduce additional dependencies on SUSY parameters. In this
analysis, the $m_{h}^{\text{max}}$ and no-mixing scenarios are studied. The corresponding excluded regions in the $\tan \beta - m_{A}$ plane are shown in Figure 3. The cross-section at each $\tan \beta - m_{A}$ point was calculated using FeynHiggs 2.5.1 by adding the $gg \rightarrow \phi$ and $b\bar{b} \rightarrow \phi$ cross-sections for a given $m_{A}$.

In the mass region $90 < m_{A} < 200$ GeV, tan $\beta$ values larger than 40-60 are excluded for the no-mixing and the $m_{h}^{\text{max}}$ benchmark scenarios. These results are the most constraining limits from the Higgs to $\tau^{+}\tau^{-}$ decay channel to date.

![Figure 2: Observed and expected 95% Confidence Level upper limit on the cross-section times branching ratio, using the neural network shown on both a log scale and a linear scale. The band represents the ±1e uncertainty on the expected limit. Also shown is the expected limit from the recent CDF result.](image)

![Figure 3: Excluded region in the $\tan \beta - m_{A}$ plane for $\mu < 0$ in a) the $m_{h}^{\text{max}}$ scenario and b) the no-mixing scenario and excluded region in the $\tan \beta - m_{A}$ plane for $\mu > 0$ in c) the $m_{h}^{\text{max}}$ scenario and d) the no-mixing scenario.](image)

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2. The ALEPH, DELPHI, L3 and OPAL Collaborations, LHFWG-Note 2005-01.
MEASUREMENT OF THE W-BOSON HELICITY FRACTIONS IN TOP-QUARK DECAYS AT CDF

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We present a measurement of the fractions \( F_0 \) and \( F_+ \) of longitudinally polarized and right-handed \( W \) bosons in top-quark decays using data collected with the CDF II detector. The data set used in the analysis corresponds to an integrated luminosity of approximately 955 pb\(^{-1}\). We select \( t\bar{t} \) candidate events with one lepton, at least four jets, and missing transverse energy. Our helicity measurement uses the decay angle \( \theta^r \), which is defined as the angle between the momentum of the charged lepton in the \( W \) boson rest-frame and the \( W \) momentum in the top-quark rest-frame. The \( \cos \theta^r \) distribution in the data is determined by full kinematic reconstruction of the \( t\bar{t} \) candidates. We find \( F_0 = 0.59 \pm 0.12 \text{(stat.)}^{+0.07}_{-0.06} \text{(syst.)} \) and \( F_+ = -0.03 \pm 0.06 \text{(stat.)}^{+0.04}_{-0.02} \text{(syst.)} \), which is consistent with the standard model prediction. We set an upper limit on the fraction of right-handed \( W \) bosons of \( F_+ \leq 0.10 \) at the 95\% confidence level.

1 Introduction

Since the discovery of the top quark in 1995 by the CDF and DØ collaborations\(^1,2\), the mass of this most massive known elementary particle has been measured with high precision. However, the measurements of other top-quark properties are still statistically limited, so the question remains whether the standard model successfully predicts these properties. In the following we present our measurement of the helicity fractions of \( W \) bosons from top-quark decay.

At the Tevatron collider, with a center-of-mass energy of 1.96 TeV, most top quarks are pair-produced via the strong interaction. In the standard model the top quark decays in nearly 100\% of all cases into a \( W \) boson and a \( b \) quark. Due to its large mass the top quark has a lifetime, that is shorter than the hadronization time. Thus, its decay products preserve the helicity content of the underlying weak interaction. The \( V - A \) structure of the weak interaction in the standard model predicts that the \( W^+ \) bosons from the top-quark decay are dominantly either longitudinally polarized or left-handed, while right-handed \( W \) bosons are heavily suppressed and
even forbidden in the limit of a massless $b$ quark. Assuming a top-quark mass of 175 GeV/c$^2$ and neglecting the mass of the $b$ quark, the fraction of longitudinally polarized $W$ bosons is predicted$^3$ to be $F_0 = 0.7$, while the fraction of left-handed $W$ bosons is $F_- = 0.3$. A significant deviation from the predicted value for $F_0$ or a nonzero value for the right-handed fraction $F_+$ could indicate new physics, such as a possible $V + A$ component in the weak interaction or other anomalous couplings at the $Wtb$ vertex.

The $W$ boson polarization manifests itself in the decay $W \to \ell \nu_\ell$ in the angle $\theta^*$, which is defined as the angle between the momentum of the charged lepton in the $W$ rest frame and the momentum of the $W$ boson in the top-quark rest-frame. The general $\cos \theta^*$ distribution is given by$^3$:

\[
\frac{dN}{d \cos \theta^*} \propto F_- \cdot \frac{3}{8} (1 - \cos \theta^*)^2 + F_0 \cdot \frac{3}{4} (1 - \cos^2 \theta^*) + F_+ \cdot \frac{3}{8} (1 + \cos \theta^*)^2
\]  

(1)

2 Event Selection and Full Reconstruction

In this analysis, we use the "lepton+jets" channel, where one $W$ boson originating from the top quarks decays leptonically into a charged lepton and a neutrino and the other $W$ boson hadronically into two quarks. Therefore, we select events with exactly one isolated electron or muon, substantial missing transverse energy due to the undetectable neutrino, and at least four jets. We require one of these jets to be tagged as $b$ jet, which means that it originates from a reconstructed secondary vertex, which is likely due to the long lifetime of $b$ hadrons. In the analyzed data set, we find 232 $t\bar{t}$ candidate events with a "lepton+jets" signature.

In order to determine the $\cos \theta^*$ distribution for the selected events, the four-vectors of the top quarks, $W$ bosons, and of the charged lepton have to be reconstructed. The full reconstruction of the entire event starts with the reconstruction of the neutrino four-vector. The transverse components are obtained from the missing transverse energy, the $z$ component is calculated using a $W$ boson mass constraint. The $W$ boson four-vector is than obtained by adding the four-vector of the neutrino and of the charged lepton. Then all possibilities to assign the jets in the event to the two $b$ quarks from the top decay and to the two quarks from the hadronic $W$ decay are considered. This treatment leads to a multiplicity of hypotheses for the reconstruction of each event. For example, an event with four jets leads to 24 possibilities for the reconstruction.

The challenge is now to find the right hypothesis for each event. Therefore we make use of constraints on the mass of the reconstructed $W$ boson, on the mass difference between the two reconstructed top quarks, which should be zero within the resolution, and on the transverse energy of the two top quarks which should in leading order be equal to the transverse energy of the entire event. Finally, we prefer hypotheses, where jets tagged as $b$ jets are assigned to the $b$ quarks from the top decays. To estimate the probability for a tagged jet to be a real $b$ jet, we use a neural network $b$-tagger. Applying this method, for each single event we choose one hypothesis from which we then obtain the $\cos \theta^*$ distribution. Figure 1 shows the distribution of the measured $\cos \theta^*$ compared to the estimated signal and background distributions.

3 Measurement

Since the number of events in the data set is small, we do not simultaneously extract the fraction of longitudinally polarized and right-handed $W$ bosons. We either fix $F_+$ to zero and fit for $F_0$, or we fix $F_0$ to its expected value and fit for $F_+$. Thus, only one free parameter is used in each fit.
To extract the single free parameter ($F_0$ or $F_+$), we use a binned maximum likelihood method. The theoretically predicted number of events in each bin is the sum of the expected background and signal. The latter is calculated from the theoretical $\cos \theta^*$ distributions (see eq. 1) for the three helicities of the $W$ bosons.

Since the reconstruction of the $t\bar{t}$ process is not perfect, acceptance and migration effects have to be considered when calculating the number of signal events ($\mu^{\text{sig,exp}}_k$) expected to be observed in a certain bin $k$ after the reconstruction:

$$\mu^{\text{sig,exp}}_k \propto \sum_i \mu^{\text{sig,thec}}_i \cdot \epsilon_i \cdot S(i, k).$$

The migration matrix element $S(i, k)$ gives the probability for an event which was generated in bin $i$ to occur in bin $k$ of the reconstructed $\cos \theta^*$ distribution. Since the event acceptance depends on $\cos \theta^*$, we weight the contribution of each bin $i$ with the efficiency $\epsilon_i$. Both $\epsilon_i$ and $S(i, k)$ are determined using the standard model Monte Carlo generator PYTHIA$^4$.

With the number of expected events and the number of observed data events in each bin, we minimize the negative logarithm of the likelihood function by varying the free parameter $F_0$ or $F_+$.

In addition, an upper limit for $F_+$ at the 95% confidence level is computed by integrating the likelihood function.

In order to compare our observations with theory, the background estimate is subtracted from the selected sample. To correct for the mentioned acceptance and reconstruction effects, we calculate a transfer function $\tau(F_0, F_+)$. The value $\tau_i$ for bin $i$ is given by the ratio of the normalized number of theoretically predicted events and the normalized number of events in this bin after applying all selection cuts and performing the reconstruction. For this calculation, we use the fit result of $F_0$ or $F_+$. Multiplying the background-subtracted number of events in bin $i$ with $\tau_i$ leads to the unfolded distribution, which then is normalized to the theoretically calculated $t\bar{t}$ production cross section $\sigma_{t\bar{t}}$ of $\sigma_{t\bar{t}} = 6.7$ pb and can directly be compared with theory-curves for different values of $F_0$ and $F_+$ (see fig. 2).

4 Results

The data used in this analysis correspond to a total integrated luminosity of 955 pb$^{-1}$, where 232 events have passed the event selection. Taking systematic uncertainties into account, assuming
a top-quark mass of 175 GeV/c^2, and assuming that the non-measured fraction is equal to its standard model expectation, the result for the fraction of longitudinally polarized and right-handed W bosons is:

\[
F_0 = 0.59 \pm 0.12 \text{ (stat) } ^{+0.07}_{-0.06} \text{ (syst)}, \\
F_+ = -0.03 \pm 0.06 \text{ (stat) } ^{+0.04}_{-0.03} \text{ (syst)}. 
\]

We obtained an upper limit on the fraction of right handed W bosons of \( F_+ \leq 0.10 \) at the 95% confidence level. Furthermore our method provides the possibility to correct the observed \( \cos \theta^* \) distribution for the selected sample for acceptance and resolution effects resulting in the distribution of the differential \( \ell \bar{\ell} \) production cross section (see figure 2). As one can see, the observation is compatible with the curves predicted by the standard model. Also the measured values for \( F_0 \) and \( F_+ \) are within the uncertainties in good agreement with the standard model predictions.

References

CONSTRATNONS ON THE EXTRA DIMENSION BY KK GRAVITINO DECAY

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We study the consequences of the gravitino decay into dark matter. We suppose that the lightest neutralino is the main component of dark matter. In our framework gravitino is heavy enough to decay before Big Bang Nucleosynthesis starts. We consider a model coming from a five dimensional supergravity compactified on $S^1/Z_2$ with gravity in the bulk and matter localized on tensionless branes at the orbifold fixed points. We require that the dark matter, which is produced thermally and in the decay of Kaluza-Klein modes of gravitino, has an abundance compatible with observation. We deduce from our model that there are curves of constraints between the size of the extra-dimension and the reheating temperature of the universe after inflation. This talk is based on hep-ph/0702183 to be published in Phys. Rev. D.

1 Introduction

The five dimensional picture of the Universe has attracted much interest in the framework of what is called brane world cosmology. In the present work we choose to work in a set-up where we assume that the radion is stabilized in a five dimensional supergravity compactified on $S^1/Z_2$ where matter and gauge fields live on the branes and gravity in the bulk ($^1$, $^2$). Low energy supersymmetry provides a natural candidate for dark matter if R-Parity is conserved and solves the hierarchy problem.

As a general framework we can choose a scenario in which susy breaking is mediated to the observable sector partly by gravity anomaly and partly by Scherk-Schwarz mechanism. It is a scenario which can avoid the appearance of tachionic masses which are present in pure anomaly mediated scenario and problems due to pure gravity mediation scenario. This mechanism provides high masses for gravitino (it means that the mass of the gravitino can be above 10 TeV).

The gravitino field has Kaluza-Klein excitations modes due to the presence of the extra-dimension. We suppose that all gravitino modes are produced after inflation during the reheating period by scattering effects in the primordial thermal bath. We suppose that the lightest mode - i.e. the zero mode - is heavy enough to mainly decay before the Big Bang Nucleosynthesis starts and we calculate its corresponding mass. Actually this is naturally the case in some scenarios of susy breaking (anomaly mediation or mix between anomaly and Scherk-Schwarz mechanism). Gravitinos modes decay into supersymmetric particles and standard model ones. If R-parity is conserved, all gravitino modes will give at the end of their decay cascade a Lightest Supersymmetric Particle (LSP) which is assumed to be the lightest neutralino. But
not all of these decay products of gravitino will contribute to the relic density of neutralinos which is assumed to be the dark matter. Actually only the gravitino modes which decay after the thermal decoupling of neutralino contribute to the dark matter density. If a gravitino mode decays before the thermal decoupling of neutralinos it does not increase the number density of neutralinos since these gravitinos produce neutralinos which are in thermal equilibrium. So a finite number of gravitinos modes contribute to the non thermally produced dark matter. The total of the thermally produced and non thermally produced amount of dark matter is constrained by the evaluation of the dark matter content of the universe. As a consequence, we can draw curves of constraints between the size of the extra-dimension and the reheating temperature because the number of KK gravitinos modes is related to the size of the extra-dimension and the number density of gravitinos is related to the reheating temperature. We have chosen a reheating temperature in the range $10^5$ GeV to $10^{10}$ GeV. This range is quite natural for scenarios which allow baryogenesis through leptogenesis. There is another constraint on the size of the extra-dimension coming from the fact that there are also KK gravitons which can disturb BBN if the number of KK modes is too high. We checked with the help of the curves given by Jedamzik\(^{17}\) that for $R^{-1} > 1$ TeV, this is not the case.

2 Interactions between gravitino and MSSM

We find (see\(^{18}\)):

$$\mathcal{L}_{\text{inter\_KK}} = \sum_{n=0}^{\infty} \frac{1}{\sqrt{2M}} e g_{ij'} \tilde{D}_\mu \phi^{i'} \chi^\alpha \sigma^\mu \sigma^\nu \psi_{n, \mu} \frac{1}{\sqrt{2M}} e g_{ij'} \tilde{D}_\nu \phi^{i'} \chi^\alpha \sigma^\mu \sigma^\nu \psi_{n, \mu} - \frac{i}{2M} e \left( \tilde{\psi}_{n, \mu} \sigma^\alpha \sigma^\mu \lambda_{(a)} + \overline{\psi}_{n, \mu} \sigma^\alpha \sigma^\mu \lambda_{(a)} \right) F^{(a)}_{\mu \lambda} \right)$$

(1)

Where $\phi$ are scalar fields, $\chi$ are chiral fermions, $\lambda$ are gauge fermions (gauginos), $F^{(a)}_{\mu \lambda}$ is the field strength tensor for the gauge boson $A^{(a)}_\mu$. Indices $i, j, \ldots$ represent species of chiral multiplets and (a), (b)\ldots are indices for adjoint representation of gauge group. $e$ is the vierbein. $g_{ij'}$ is the Kahler metric and $\sigma_{n, \mu}$ is the KK gravitino field for the $n$th mode.

The masses of the modes $n$ are related to the mass of the zero mode by the relation\(^{12}\):

$$M_n = M_0 + \frac{n}{R}$$

(2)

3 Abundances and lifetime

Abundance for the zero mode is given by Kohri and al\(^{18}\) taking into account production during the inflaton-dominated epoch and Fradler and al\(^{16}\) have the same result without taking into account production during inflation. This result is given for masses of gravitino much higher than gauginos masses but the calculation of the creation term in the Boltzmann equation is made with particles without masses: their mass is supposed negligible compared to the average energy in the center of mass of each reaction. The formula is:

$$Y_{3/2} \simeq 1.9 \times 10^{-12} \times \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \left[ 1 + 0.045 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right] \left[ 1 - 0.028 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right],$$

(3)

Where $Y_{3/2} = \frac{n_{3/2}}{s}$, $n_{3/2}$ is the number density, $s$ is the entropy density and $T_R$ the reheating temperature. The quantity $Y = \frac{n}{s}$ is the density per comoving volume. We have to take into account in our calculation the different masses of gravitino modes. We find as a good approximation this rule for the abundance of the different modes:

$$Y_{3/2}^k = Y_{3/2}^0, \quad \text{for } M^k \leq T_R \text{ and }$$

$$Y_{3/2}^k = 0, \quad \text{for } M^k > T_R,$$

(4)
Where $k$ represents the index of the KK mode, $M^k$ is the mass of the $k^{th}$ gravitino mode and $Y^k_{3/2}$ its abundance.

We have calculated the lifetime for heavy gravitino (masses $> 10$ TeV) using ($3'$):

$$
\tau_k = 1.4 \times 10^7 \times \left( \frac{M_k}{100 \text{ GeV}} \right)^{-3} \text{ Sec}
$$

(5)

### 4 Neutralinos and equation of constraints

In our model the LSP is the lightest neutralino. We choose to work with a mass of LSP equal to 120 GeV. This analysis can be easily rescaled for another choice for the mass: we also show results for a mass of the LSP equal to 200 GeV in the article$^6$. The dark matter density is:

$$
0.106 < \Omega \ h^2 < 0.123
$$

(6)

We call $\Omega_{\chi h}$ the thermal density of neutralinos. We find this approximate relation between $\Omega_{\chi h}$ and $x_f$:

$$
\Omega_{\chi h} \ h^2 = 3.61 \times 10^6 \frac{m_{\text{LSP}}}{1 \text{ GeV}} \ x_f^2 \ e^{-x_f}
$$

(7)

We choose different values of $\Omega_{\chi h} \ h^2$ and complete with the non thermal production coming from the gravitino decay. We call this non thermal production $\Delta \Omega \ h^2$.

$$
0.106 \leq \Omega_{\chi h} \ h^2 + \Delta \Omega \ h^2 \leq 0.123
$$

(8)

As one gravitino produces one neutralino, we can write:

$$
\Delta \Omega \ h^2 = \frac{m_{\text{LSP}}}{\rho_c} \ h^2 \sum_{k=0}^{k=n} Y^k_{3/2}
$$

(9)

The index $n$ corresponds to the last mode to be taken into account. It is the mode decaying just when the LSP decouples from the thermal bath. Only the gravitino modes decaying after the thermal decoupling of the neutralinos contribute to increase the quantity of neutralinos.

Using the equations (3) and (4), and the equation (9) we find a constraint equation coming from the relation (8):

$$
R^{-1} = (M^\tilde{g} - M^G) \times
$$

$$
\left( \frac{\Delta \Omega \ h^2 \rho_c}{m_{\text{LSP}} \ s_0 \ h^2} \ 1.9 \times 10^{-12} \times \frac{\Omega_{\chi h}}{M_{\tilde{g}}} \ [1 + 0.045 \ln \left( \frac{M_{\tilde{g}}}{112 \text{ GeV}} \right) \ [1 - 0.028 \ln \left( \frac{M_{\tilde{g}}}{112 \text{ GeV}} \right)] - 1 \right)^{-1}
$$

(10)

### 5 Results

In the paper$^6$, we have treated three cases for different values of $\Omega_{\chi h} \ h^2$. As an example we present below two figures of the case where $\Omega_{\chi h} \ h^2 = 0.053$, $\Delta \Omega_{\chi h} \ h^2_{\text{min}} = 0.033$, $\Delta \Omega_{\chi h} \ h^2_{\text{max}} = 0.070$, $M_{\tilde{g}} = 9.66 \times 10^6 \text{ GeV}$.

### 6 Conclusion

The results that we obtain are independent from the sasy mass spectrum since the gravitino is heavy enough to make negligible the influence of other susy particles. The results show that the size of the radius $R$ is not only bounded by a maximum value but also by a minimum value in a wide range of possible values for the thermal production of neutralinos and in a wide range of values for the reheating temperature. To obtain a minimum size for the radius is a new result.

Kaluzza-Klein modes of graviton are also present and may disturb BBN. We checked with an approximate method that for $R^{-1}$ above 1 TeV it is not the case. This already implies a bound on the reheating temperature which can not be lower than a minimum value in the cases where the radius $R$ is bounded by a minimum value.
Figure 1: $T_R$ less than $9.66 \times 10^6$ GeV. Only the band between the two diagonal curves is allowed. The zone below the straight line $R^{-1} = 1$ TeV is excluded if KK gravitons constraint is taken into account.

Figure 2: $9.66 \times 10^6$ GeV $\leq T_R \leq 10^8$ GeV. Only the band between the two curves is allowed.

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References

Evidence For Single Top Quark Production Using The Matrix Element Analysis Technique in 1 fb$^{-1}$ of Tevatron RunII Data

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We present the first evidence for electroweak single top quark production using nearly 1 fb$^{-1}$ of Tevatron Run II data at $\sqrt{s} = 1.96$ TeV. We select single-top-like data events in the lepton+jets decay channel and isolate them from backgrounds using the matrix element analysis method. This technique employs leading order matrix elements to compute an event probability for both signal and background hypotheses. Using the expected signal acceptance, background, and observed data we measure the single top quark cross section:

$$\sigma(p\bar{p} \rightarrow tb + tq\bar{b} + X) = 4.6^{+1.8}_{-1.5} \text{ pb}$$

The probability for the background to have fluctuated up to give at least the cross section measured in this analysis is 0.21%, which corresponds to a Gaussian equivalent significance of 2.9$\sigma$.

1 Introduction

The top quark is the heaviest of the known fermions and is currently only produced at Fermilab’s Tevatron proton-antiproton collider with a center-of-mass energy of 1.96 TeV. While most top quarks are produced via the well measured QCD pair production mode ($p\bar{p} \rightarrow tt\bar{t}$), they can also be created via an electroweak interaction known as single top. At the Tevatron, single top quarks are produced in two modes: the s-channel process ($p\bar{p} \rightarrow tb$) which has an estimated cross section of $0.88 \pm 0.11 \text{ pb}$ and the t-channel process ($p\bar{p} \rightarrow tq\bar{b}$) with a cross section of $1.98 \pm 0.25 \text{ pb}$. Feynman diagrams for the two production channels are shown in Figure 1.

Measuring single top quark production is interesting because one can directly determine the magnitude of the CKM matrix element $V_{tb}$ since $\sigma_{tb+tq\bar{b}} \propto |V_{tb}|^2$. Single top quark production is also sensitive to new physics. The existence of a new charged gauge boson ($W'$) would enhance
the effective $s$-channel cross section, while flavor changing neutral currents in the top sector would enhance the effective $t$-channel cross section.

2 Event Selection and Background Modeling

We measure single top quark production in the lepton+jets decay channel. These events are selected in data by requiring the following selection cuts.

- One high $p_T$ isolated lepton (electron or muon). The electron is required to be measured in the central calorimeter ($|\eta^{\text{jet}}| < 1.1$) with $p_T > 15$ GeV. The muon must have $p_T > 18$ GeV and $|\eta^{\text{jet}}| < 2.0$. Events with electrons veto events with muons and visa versa.

- Missing transverse energy > 15 GeV.

- Between two and three reconstructed jets. The leading jet (highest $p_T$) must have $p_T > 25$ GeV and $|\eta^{\text{jet}}| < 2.5$. The second jet is required to have $p_T > 20$ GeV and the third jet must have $p_T > 15$ GeV. Both the second and third jets much have $|\eta^{\text{jet}}| < 3.4$.

- At least one jet must be $b$-tagged by the neural network $b$-tagger. This analysis selects events with one or two $b$-tagged jets.

After applying this event selection there are three main backgrounds present in the data: $W$+jets, $t\bar{t}$, and QCD multijet production. The $W$+jets and $t\bar{t}$ backgrounds are modeled using ALPGEN with the MLM jet-parton matching scheme. The shape for QCD multijet events is derived from data events that pass all selection criteria except lepton-jet isolation. The $t\bar{t}$ background is normalized to the NLO Tevatron cross section, while $W$+jets and QCD multijet events are normalized to the data before $b$-tagging.

Both $s$-channel and $t$-channel single top quark events are modeled using the SINGLETOP Monte Carlo generator. A table showing the expected number of single and background events can be found on the DØ single top public webpage.

3 Matrix Element Analysis Technique

The matrix element analysis technique attempts to assign events as signal or background-like based on the normalized proton-antiproton differential cross section as shown in Equation 1

$$ P(\vec{x}) = \frac{1}{\sigma} \frac{d\sigma}{d\vec{x}} = \frac{1}{\sigma} \sum_{i,j} \int dy dy' \left[ f_i(q_1,Q^2) dq_1 \times f_j(q_2,Q^2) dq_2 \times \frac{d\sigma_{HH,ij}}{dy} \times W(\vec{x},\vec{y}) \right] ; \quad (1) $$

where $f_i(q_1,Q^2)$ is the parton distribution function for parton flavor $i$, $d\sigma_{HH,ij}/dy$ is the parton-parton hard scatter differential cross section ($\propto |\!| M|^2 |\!|$), $\sigma$ is the normalization factor, $\int dy dq_1 dq_2$ represents the integration over the parton phase space, and $W(\vec{x},\vec{y})$ is the detector resolution function that maps the parton state ($\vec{y}$) to the detector state ($\vec{x}$).
The probability density in Equation 1 is evaluated separately for events with two and three jets. Events with two jets are evaluated using two single top matrix elements and three background $W+\text{jets}$ matrix elements as shown in Figure 2.

![Figure 2: Leading order Feynman diagrams used for events with two jets. From left to right: $p\bar{p} \rightarrow tb$ (signal), $p\bar{p} \rightarrow t\bar{t}q$ (signal), $p\bar{p} \rightarrow Wb\bar{b}$ (background), $p\bar{p} \rightarrow Wc\bar{c}$ (background), and $p\bar{p} \rightarrow Wg\gamma$ (background).](image)

Events with three jets are evaluated using two single top matrix elements and one background $W+\text{jets}$ matrix element as shown in Figure 3.

![Figure 3: Leading order Feynman diagrams used for events with two jets. From left to right: $p\bar{p} \rightarrow tbg$ (signal), $p\bar{p} \rightarrow t\bar{q}b$ (signal), and $p\bar{p} \rightarrow Wbb$ (background).](image)

The signal and background probabilities are combined using the \textit{a-posteriori} Bayesian probability density for the signal process (either $s$-channel or $t$-channel) given the event $\vec{x}$ as shown in Equation 2

$$D(\vec{x}) = \frac{P_{\text{Signal}}(\vec{x})}{P_{\text{Signal}}(\vec{x}) + P_{\text{Background}}(\vec{x})}. \quad (2)$$

The discriminant variable defined in Equation 2 is evaluated for all data, signal Monte Carlo, and background events. Figure 4 shows the expected distributions of $s$-channel, $t$-channel, and $Wbb$ Monte Carlo events in a two-dimensional plane defined by the $t$-channel and $s$-channel discriminant.

![Figure 4: Expected performance of $s$-channel (left), $t$-channel (middle), and $Wbb$ (right) Monte Carlo events shown in a two-dimensional plane defined by the $t$-channel and $s$-channel discriminant.](image)

4 Results with 0.9 fb$^{-1}$

Applying the matrix element discriminant to signal, background, and data event resulted in an excess of data events in the signal region. From this excess, we measure the single top cross section as

$$\sigma(p\bar{p} \rightarrow tb + t\bar{q}b + X) = 4.6^{+1.8}_{-1.5} \text{ pb.}$$

The full matrix element discriminant including a zoom of the signal region is shown in Figure 5.

The significance of the observed result is determined using ensemble tests. In this test, pseudo datasets with zero signal content are created and the number of datasets that yield a
cross section greater than or above the measured cross section is determined. Using nearly 100,000 ensemble datasets, this fraction was measured to be 0.21% corresponding to a 2.9σ signal significance. The agreement of this result with the standard Model was measured using pseudo datasets with the signal fraction set to the standard model value. In this test, 21% of the datasets resulted in a cross section greater than or equal to the measured value. Information regarding the systematics errors in this analysis as well as the Bayesian statistical technique used in the cross section measurement can be found in our recently published article in Physical Review Letters.8

5 Summary

We presented evidence for electroweak single top quark production at the Tevatron. An analysis of nearly 1 fb−1 of Run II data using the matrix element method to select single top quark-like events measures the combined s + t-channel production cross section to be

\[ \sigma(p\bar{p} \rightarrow tb + t\bar{q}b + X) = 4.6^{+1.8}_{-1.5} \text{ pb}, \]

where the probability of a background fluctuation is 0.21%, which corresponds to a Gaussian equivalent signal significance of 2.9σ.

6 Acknowledgments

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SEARCH FOR SQUARKS AND GLUINOS AT DØ

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A search for scalar quarks and gluinos is performed with 0.96 fb⁻¹ of data collected by the DØ experiment in pp collisions at $\sqrt{s} = 1.96$ TeV at the Fermilab Tevatron Collider. The topologies analyzed consist of acoplanar jets and multijet events with missing transverse energy. We find the data to be consistent with Standard Model expectations, and set 95% C.L. exclusion domains in the squark and gluino mass plane and in the $(m_0, m_{1/2})$ plane, within the framework of minimal supergravity with $\tan\beta = 3$, $A_0 = 0$ and $\mu < 0$.

1 Introduction

Supersymmetric theories predict for each elementary fermion (resp. boson) a new elementary boson (resp. fermion). These new particles have the same quantum numbers as their Standard Model (SM) partners, but a spin which differs by one half. Particles carry a new multiplicative quantum number $R$, which is 1 for SM particles and $-1$ for super-particles. The partners of quarks (resp. gluons) are called squarks (resp. gluinos) and have a spin 0 (resp. 1/2).

In R-parity conserving theories, supersymmetric particles are thus produced in pairs and final states of their decay chains contain the lightest supersymmetric particles (LSP) which is stable. The supersymmetric partners of neutral gauge and Higgs bosons are called neutralinos, and the lightest one, $\tilde{\chi}_1^0$, arises as the natural LSP in supergravity inspired models\(^1\).

The generic models predict squarks of the four lightest quark flavors to have similar masses, and this, independently of the helicity of the quark. If they are sufficiently light, squarks can be widely produced at Tevatron and they decay according to $\tilde{g} \to q\tilde{\chi}_1^0$ as shown in Fig. 1.a. Tevatron can also produce a large amount of gluinos if they are not too heavy. In particular, if they are lighter than squarks, their preferred decay is then $\tilde{g} \to q\bar{q}\tilde{\chi}_1^0$ (Fig. 1.b).

The $\tilde{\chi}_1^0$ not being detected, final state topologies for such events are two jets and missing transverse energy ($E_T$) in the case of squark pair production (Fig. 2.a), at least three jets and
\( \mathcal{E}_T \) for simultaneous production of squark and gluino (Fig. 2.b), and four jets and \( \mathcal{E}_T \) for pair production of gluinos (Fig. 2.c).

This search is performed in the minimal supergravity (mSUGRA) framework\(^1\), with the following parameters fixed: \( \tan \beta = 3 \), \( A_0 = 0 \) and \( \mu < 0 \). In order to increase the sensitivity, three analyses are developed, searching events with two (dijet analysis), three (3-jets analysis) or at least four (gluino analysis) jets and \( \mathcal{E}_T \).

2 Common part of the three analyses

The three analyses start with a common data set and preselection which are later on optimized. The data sample used corresponds to an integrated luminosity of 0.96 fb\(^{-1}\). At the trigger level it selects events with 2 acoplanar jets and \( \mathcal{E}_T \), or with multi-jets and \( \mathcal{E}_T \).

The main backgrounds are multijet events with \( \mathcal{E}_T \) coming from mismeasured jets (instrumental background), and "electroweak" background: \( Z^0(\nu \nu) + \) jets, \( W^\pm(\ell v) + \) jets (with \( l \) a charged lepton) and \( t\bar{t} \). Because the signals have jets with high transverse energy (\( \mathcal{E}_T \)), one selects events with at least two jets with \( \mathcal{E}_T > 35 \) GeV in the central region of the calorimeter (|\( \eta \)| < 0.8, where \( \eta = \ln(\tan(\theta/2)) \), and \( \theta \) is the polar angle relative to proton beam). In order to reduce instrumental background, jets are track confirmed and the two leading jets must be acoplanar: \( \Delta \phi(jet_1, jet_2) < 165^\circ \) (\( \phi \) being the azimuthal angle). The instrumental background is greatly reduced by requiring \( \mathcal{E}_T > 75 \) GeV (Fig. 3.a) and an isolation of \( \mathcal{E}_T \). For example, the minimal angle between \( \mathcal{E}_T \) and jets has to be greater than 40° for the dijet analysis (Fig. 3.b). Finally veto on isolated muon and electron is performed to fight against electroweak background.

3 Optimization and combination of the three analyses

After preselection the three analyses are optimized by adjusting cuts on \( \mathcal{E}_T \) and \( \mathcal{H}_T \) which is the scalar sum of the \( \mathcal{E}_T \) of the jets. The optimized values of \( \mathcal{E}_T \) and \( \mathcal{H}_T \) and listed in Tab. 1.a. No excess of events is found, data being in agreement with SM expectation for all analyses (Tab. 1.b).

In order to increase the overall sensitivity, the three optimized analyses are then combined, giving rise to seven independent samples. Systematic uncertainties are re-computed for each combination, and a global limit is determined using the modified frequentist approach\(^2\) including...
Figure 3: (a) Distribution in $E_T$ before the two final cuts on $H_T$ and $E_T$ in the dijet analysis. The signal $M_q = 375$ GeV and $M_l = 416$ GeV is drawn as a hatched histogram. (b) Distribution in $\Delta\phi_{\text{min}}(E_T\text{, any jet})$ in the dijet analysis before cutting on that variable.

### Table 1: Results of the optimization.

<table>
<thead>
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<th>Analysis</th>
<th>$E_T$ (GeV)</th>
<th>$H_T$ (GeV)</th>
</tr>
</thead>
<tbody>
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<td>dijet</td>
<td>225</td>
<td>300</td>
</tr>
<tr>
<td>3-jets</td>
<td>150</td>
<td>400</td>
</tr>
<tr>
<td>gluino</td>
<td>100</td>
<td>300</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Analysis</th>
<th>Data</th>
<th>Total Background</th>
</tr>
</thead>
<tbody>
<tr>
<td>dijet</td>
<td>5.0</td>
<td>7.5 ± 1.1 (stat.) +1.3 (syst.)</td>
</tr>
<tr>
<td>3-jets</td>
<td>6.1</td>
<td>6.1 ± 0.4 (stat.) +1.2 (syst.)</td>
</tr>
<tr>
<td>gluino</td>
<td>33.4</td>
<td>33.4 ± 0.8 (stat.) +1.6 (syst.)</td>
</tr>
</tbody>
</table>

4 Final Results

The main systematic uncertainties of this search are the background cross-sections (15%), the jet energy scale (from 6 to 17%), the luminosity calculation (6.1%), the effect of parton distribution functions (PDF) uncertainties on signal acceptance (6%) and track confirmation (5%).

The signal cross-section uncertainties are estimated using 41 CTEQ6.1M PDF sets. The PDF effect is combined quadratically with the effect of the renormalization and factorization scale by a factor of two up or down. The three resulting limits are presented in yellow and red color in the squark and gluino mass plane (Fig. 4), and in the $(m_0, m_{1/2})$ plane (Fig. 5). The previous limits from Tevatron and LEP are significantly improved. Using the most conservative excluded cross-sections, gluino and squark masses are excluded up to 289 and 375 GeV respectively at 95% C.L.

References

Figure 4: In the squark and gluino mass plane, regions excluded by the analysis at 95% C.L. in the mSUGRA framework for $\tan\beta = 3$, $A_0=0$ and $\mu < 0$. The red line is the excluded region for the central PDF and renormalization and factorization scale ($\mu_{r,f} = Q$). The yellow band is obtained by varying $\mu_{r,f}$ by a factor of two and combining with the effects of the PDF uncertainties.

Figure 5: In the $(m_0, m_{1/2})$ plane, regions excluded by the analysis at 95% C.L. in the mSUGRA framework for $\tan\beta = 3$, $A_0=0$ and $\mu < 0$. The red line is the excluded region for the central PDF and renormalization and factorization scale ($\mu_{r,f} = Q$). The yellow band is obtained by varying $\mu_{r,f}$ by a factor of two and combining with the effects of the PDF uncertainties. There is no mSUGRA solution in the dark grey region. LEP2 chargino and slepton searches excluded the beige and the green regions respectively. The nearly horizontal thin black lines are the gluino iso-mass curves corresponding to gluino masses of 150, 300, 450 and 600 GeV. The other ones are squark iso-mass curves corresponding to squark masses of 150, 300, 450, 600 and 750 GeV.
THE COMMISSIONING OF THE ATLAS CALORIMETERS WITH COSMIC MUONS

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The commissioning of the ATLAS calorimeters is an ongoing process since early 2006. During this period, cosmic muons have been recorded in several runs combining both hadronic and electromagnetic calorimeters. Among the goals are the measurement of the uniformity of the liquid argon electromagnetic calorimeter to the level of 1% and the intercalibration in time of its channels to 1 ns.

1 Commissioning with cosmics

The ATLAS EM calorimeters undergo in 2007 the last steps of their installation. During the previous two years of commissioning, several milestones have been successfully attained, most prominently the complete cabling of the whole electronics chain, the cryostats fill with liquid argon and the high voltage ramping. The calorimeters in this definitive setup could then detect the cosmic muons reaching them in spite of their location 150 m underground in the cavern.

1.1 ATLAS calorimetry during the commissioning

The commissioning involves all calorimeter subdetectors for which a detailed description can be found in the ATLAS Technical Design Report\textsuperscript{1}. The endcap regions being operational only since a few months, the emphasis will be put on the barrel part from which more information can be extracted with this special data taking. The trigger signal is provided by the hadronic calorimeter detecting top-bottom coincidences at a rate of one event every 20 seconds. Monte-carlo simulations had predicted a strong occurrence of incoming muons making their way down the access shafts. The accumulated data corroborate this idea with a maximum of events originating...
from that direction at $|\eta| \sim 0.3$. The coverage of the trigger ranges basically up to $|\eta| < 0.8$ as seen in fig.1, the rates falling abruptly beyond this rapidity.

The liquid argon technology used as active material of the electromagnetic calorimeter was adopted for its high stability and radiation hardness. This sensitive part is located in the gaps between the absorbers, accordion shaped plates made of lead. The schema of fig.1 shows the longitudinal segmentation of this detector in a front, middle and back compartments. Due to its larger depth, the middle layer collects most of the small signals left by muons and will hence be subject to a more detailed analysis.

1.2 Motivations

Since the ATLAS Collaboration unveiled the first data taking of cosmic muons in August 2006 (fig.2), an effort has been initiated in order to improve the knowledge of the calorimeters. Though the preceding testbeam data have brought much to our understanding, it is an opportunity to assess the behaviour of still untested modules.

Muons propagating almost at the minimum of ionisation do not offer optimal conditions for measuring performances of the EM calorimeter. The small energy deposits they produce amount typically to 275 MeV in the middle, only a factor 8 above the electronic noise and two orders of magnitude below the energy scale relevant to the LHC physics. Unlike the LHC environment, cosmic muons produce asynchronous data and their trajectories are seldomly projective, i.e. passing close to the interaction point. On the other hand, muons have well-understood energy deposition mechanisms and, unaffected by non-uniformities in the absorbers, they give a direct sensitivity to the width of the liquid argon gaps and effects due to a wrong calibration or signal reconstruction. They also allow to scan a large fraction of the detector without concerns about the energy of the incoming particles. A simulation of projective muons has demonstrated that a precision of 1% on the value of the most probable energy deposit could be achieved for each $\eta$ range of the middle with the statistics accumulated during 3 months.

Another motivation comes from the observation of large energy deposits in the calorimeter interpreted as muons emitting a photon which then produces an electromagnetic shower. In the figure 2, an example of a high amplitude signal is shown, representing an equivalent of 25 GeV. These events can be exploited in various ways, for instance to validate predictions of the physics pulse shapes or to be used as reference when intercalibrating in time the channels. From the
current observed rate, we can expect at least one such event per cell after the planned data taking period (3 months). Reversely, the absence of signals identifies dead channels. An amount of 60k muon signals have been recorded during the 3 days long runs taken in fall 2006 with a limited coverage and a high voltage set to 1600V. In the first half of 2007, more regular runs were taken overnight and during weekends with a more complete setup.

2 Results

The challenges related to the analysis of cosmic muons had to be overcome. For instance, the optimal filtering method\(^2\) used for the signal reconstruction in the EM calorimeter does require the knowledge of the pulse’s time. Though it can be evaluated \textit{in situ} with large signals, this information will be provided by the hadronic calorimeter. Another issue concerns the signal to noise ratio which should be improved as much as possible. The low event rate allows data taking with more samples than the nominal 5 at the LHC. A significant reduction of the noise is obtained when performing instead an optimal filtering with 29 samples (see fig.3). Is also shown the Landau distribution of the deposited energy of the muons.

Figure 3: The noise is reduced by a factor 1.8 if the number of samples for the optimal filtering is increased from 5 to 29 (left). The graph on the right is the Landau distribution of muon signals observed during the testbeam. The electronic noise is taken into account by a convolution with a gaussian.
2.1 Uniformity in energy

Due to the accordion geometry, a signal is always shared between two neighbouring cells in $\phi$. Adding these two cells would contain the whole signal as long as the muon is projective enough to not leak in adjacent cells in $\eta$ for instance. Criteria to assess the biases introduced by the non-projectivity are under study. At this point, signals in the front could be used to determine more precisely the paths and hence avoid this issue. Out of the accumulated data, around 5% are described as projective enough to be considered in a uniformity measurement. Though each muon produces signals at the top and bottom of the calorimeter, the limited statistics of projective muons allows to seek for effects along $\eta$ (fig.4) only with the assumption of symmetries of the detector in $\phi$, top versus bottom and between A or C sides.

![Graph showing comparison between data and MC of the peak of the Landau shape for various $\eta$ ranges. The time resolution for muons as a function of the energy is presented on the right.](image)

Figure 4: On the left, this preliminary graph shows a comparison between data and MC of the peak of the Landau shape for various $\eta$ ranges. The time resolution for muons as a function of the energy is presented on the right.

2.2 Intercalibration in time

To perform an intercalibration in time of two channels, a single particle crossing both of them would be sufficient after a correction for the time-of-flight. Unfortunately, the time measurement remains very unprecise at low amplitudes (see fig.4). Large signals allow to measure precisely the relative offsets between the tile and electromagnetic signals. With more data being analysed, these reference signals will help to disentangle the various sources of offsets and solve the system such that all the channels are intercalibrated to 1 ns.

3 Conclusion

With more components becoming operational in ATLAS, recent cosmic runs taken in Spring 2007 are already integrating parts of the inner detector (TRT) and the muon system. Uniformity analysis and time intercalibration are currently performed in order to understand as much as possible the calorimeters before the LHC start-up next year.

References

Young Scientist Forum

2 - Flavour Physics
Charged kaon lifetime at KLOE

THE KLOE COLLABORATION
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Preliminary result on the charged kaon lifetime $\tau^+$, obtained by the KLOE experiment operating at DAΦNE, the Frascati $\phi$-factory, is presented.

1 DAΦNE and KLOE

The DAΦNE $e^+e^-$ collider operates at a total energy $\sqrt{s} = 1020$ MeV, the mass of the $\phi(1020)$-meson.

Since 2001, KLOE has collected an integrated luminosity of about 2.5 $fb^{-1}$. Results presented below are based on 2001-02 data for about 450 $pb^{-1}$.

The KLOE detector consists of a large cylindrical drift chamber, DC, surrounded by a lead/scintillating-fiber electromagnetic calorimeter, EMC. The drift chamber $^1$, is 4 m in diameter and 3.3 m long. The momentum resolution is $\sigma(p_T)/p_T \sim 0.4\%$. Two track vertices are reconstructed with a spatial resolution of $\sim 3$ mm. The calorimeter $^2$, composed of a barrel and two endcaps, covers 98% of the solid angle. Energy and time resolution are

\( \sigma(E)/E = 3.7\%/\sqrt{E[GeV]} \) and \( \sigma(i) = 57ps/\sqrt{E[GeV]} \pm 100ps \). A superconducting coil around the detector provides a 0.52 T magnetic field.

The KLOE trigger\(^4\), uses calorimeter and drift chamber information. For the present analysis only the calorimeter triggers have been used. Two energy deposits above threshold, \( E > 50 \) MeV for the barrel and \( E > 150 \) MeV for the endcaps, have been required.

2 The tag mechanism

In the center of mass the \( \phi \)-meson decays into anti-collinear \( K\bar{K} \) pairs. In the laboratory this remains approximately true because of the small crossing angle of the \( e^+e^- \) beams. Therefore the detection of a \( K\bar{K} \) tags the presence of a \( K\bar{K} \) of given momentum and direction.

The decay products of the \( K^+ \) pair define two spatially well separated regions called the tag and the signal hemisphere. Identified \( K^+ \) decays tag a \( K^\mp \) beam and provide an absolute count, using the total number of tags as normalization. This procedure is a unique feature of a \( \phi \)-factory and provides the means for measuring absolute branching ratios.

Charged kaons are tagged using the two body decays \( K^+ \rightarrow \mu^+\nu_\mu \) and \( K^- \rightarrow \pi^\pm\pi^0 \). Since the two body decays correspond to about 85% of the charged kaon decays\(^4\) and since \( BR(\phi \rightarrow K^+K^-) \approx 49% \), there are about \( 1.5 \times 10^6 K^+K^- \) events/pb\(^{-1}\).

The two body decays are identified as peaks in the momentum spectrum of the secondary tracks in the kaon rest frame and using the pion mass hypothesis \( p^m \) (Fig. 1). In order to minimise the impact of the trigger efficiency, the tags must provide the Euc trigger of the event, so called self-triggering tags. \( N_{\text{self tagging}} \approx 2 \times 10^5 \) per pb\(^{-1}\).

![Momentum spectrum in the kaon rest frame of the negative charged decay particle, assuming the pion mass for data (dots) and MC (lines). The two peaks correspond to pions and muons from \( K^- \rightarrow \pi^-\pi^0 \) (205 MeV/c) and \( K^- \rightarrow \mu^-\nu_\mu \) (230 MeV/c).](image)

3 Measurement of the charged kaon lifetime

The measurement is performed using 230 pb\(^{-1}\) collected at \( \phi \) peak. The data sample has been split in two uncorrelated subsamples, 150 pb\(^{-1}\) have been used for the measurement, the remaining 80 pb\(^{-1}\) have been used to evaluate the efficiencies. Both charged \( K_{\mu2} \) tags have been used.

There are two methods available for the measurement: the kaon decay length and the kaon decay time. The two methods allow cross checks and studies of systematics; their resolutions are comparable.
The method relying on the measurement of the charged kaon decay length requires first the reconstruction of the kaon decay vertex in the fiducial volume using only DC information: the signal is given by a $K^\pm$, moving outwards in the DC with momentum $70 < p_K < 130$ MeV/c; and having point of closest approach with $0 < \sqrt{x_{PC: A}^2 + y_{PC: A}^2} < 10$ cm and $|z_{PC: A}| < 20$ cm.

The kaon decay vertex in the DC fiducial volume ($40 < \sqrt{x_{V}^2 + y_{V}^2} < 150$ cm, $|z_{V}| < 150$ cm) is required. Once the decay vertex has been identified the kaon track is extrapolated backward to the interaction point into 2 mm steps, taking into account the ionization $dE/dx$ to evaluate its velocity $\beta c$.

Then the proper time is obtained from the equation:

$$\tau^* = \sum_i \Delta T_i = \sum_i \frac{\sqrt{1 - \beta_i^2}}{\beta_i} \Delta l_i$$

(1)

The efficiency has been evaluated directly from data. The control sample has been selected using calorimetric information only, selecting for a neutral vertex: two clusters in time fired by the photons coming from the $\pi^0$ decay.

The proper time is fitted between 16 and 30 ns correcting for the efficiency. Resolution effects have been taken into account.

The preliminary result we have obtained for the $K^+$ is:

$$\tau^+ = (12.367 \pm 0.044 \pm 0.065) \text{ ns}$$

(2)

with $\chi^2 = 17.7/15$, corresponding to a $\chi^2$ probability $P(\chi^2) = 28.4\%$. The evaluation of systematic uncertainties is still preliminary, final study will be presented at the conference.

![Graph](image)

Fig. 2: Charged kaon proper time distribution, obtained with the first method, fitted (red line) with a convolution of an exponential and a resolution function.

The second method relies on the measurement of the kaon decay time. It requires the backward extrapolation to the interaction point of the tagging kaon track and the forward
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3 - Astrophysics and Cosmology
ANALYSIS OF THE SN1987A TWO-STAGE EXPLOSION HYPOTHESIS
WITH ACCOUNT FOR THE MSW NEUTRINO FLAVOUR CONVERSION

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117218, B.Cheremushkinskaya 25, Moscow, Russia

Detection of 5 events by the Liquid Scintillation Detector (LSD) on February, 23, 1987 was interpreted in the literature as a detection of neutrinos from the first stage of the two-stage supernova collapse. We pose rigid constraints on the properties of the first stage of the collapse, taking into account MSW neutrino flavour conversion and general properties of supernova neutrino emission. The constraints depend on the unknown neutrino mass hierarchy and mixing angle $\theta_{13}$.

SN1987A was the only supernova to date which produced a measured neutrino signal. Four experiments reported the detection of neutrinos: LSD $^{2,2}$ (Liquid Scintillator Detector), KII $^{2}$ (Kamiokande II), IMB $^{7}$ (Irvine-Michigan-Brookheaven) and BST $^{7}$ (Baksan Scintillator Telescope). While LSD registered neutrino burst at 2:52UT, February 23, the other three experiments – at 7:35UT February 23 (UT stands for Unitary Time). Each experiment reported only one burst: LSD observed no statistically significant counterpart for the KII, IMB and BST neutrino signals and vice versa. This puzzling discrepancy could be, in principle, explained by a two-stage supernova collapse hypothesis, as was stated in a number of works. $^{2,2,7,7}\text{.}$ Various two-stage supernova collapse scenarios can be posed, MSW effect $^{7,7}$ in the matter of the star being of crucial importance. For example, it was shown $^{7}$ that accounting for the neutrino flavour conversion can spoil the reported $^{7}$ concordance of the rotating collapsar model with the data. Analysis, independent of the particular collapse model, was also performed. $^{7}$ In the present note we report the results of the extended analysis. $^{7}$ It provides a more elaborate statistical study of the data and accounts for the supernova shock wave effect, which may influence neutrino flavour conversion. $^{7}$
Table 1: Type, working material and working mass (in tons) of the detectors, numbers of events \( (N_{\nu}) \) at 2:52 and 7:35 (according to the cited references).

<table>
<thead>
<tr>
<th>Type</th>
<th>LSD</th>
<th>KII</th>
<th>IMB</th>
<th>Baksan</th>
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</thead>
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<tr>
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<td>scintillator</td>
<td>cherenkov</td>
<td>cherenkov</td>
<td>scintillator</td>
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<tr>
<td>( C_{n}H_{2n} )</td>
<td>( Fe )</td>
<td>90t</td>
<td>200t</td>
<td>( H_{2}O )</td>
</tr>
<tr>
<td>( N_{\nu} ) at 2:52</td>
<td>5 (^{+1}_{-2})</td>
<td>2 (^{+2}_{-1})</td>
<td>0 (^{+1}_{-2})</td>
<td>1 (^{+1}_{-2})</td>
</tr>
<tr>
<td>( N_{\nu} ) at 7:35</td>
<td>2 (^{+2}_{-1})</td>
<td>11 (^{+2}_{-1})</td>
<td>8 (^{+2}_{-1})</td>
<td>6 (^{+2}_{-1})</td>
</tr>
</tbody>
</table>

Detector characteristics and numbers of registered events at 2:52 and 7:35 are given in Table 1. We discuss only the first stage of the presumable two-stage collapse and, accordingly, only the first neutrino signal, which occurred at 2:52. Moreover, we compare only LSD and KII signals, and do not use IMB and BST data.

It should be noted that the imitation rate from the background for the LSD event cluster at 2:52 was fairly small - 0.7 per year.\(^{7,7}\) This justifies the attempts to find an explanation for the LSD neutrino signal.

Reactions essentially relevant for neutrino and antineutrino detection in LSD and KII are listed below. Each reaction is relevant for the detector(s) indicated in parentheses.

\[
\bar{\nu}_e + p \rightarrow e^+ + n \quad \text{(KII, LSD)}
\]

\[
\begin{align*}
\nu_e + Fe & \rightarrow e + Co^* \quad \text{(LSD)} \\
\nu_\mu + Fe & \rightarrow \mu + Fe^* \quad \text{(LSD)} \\
\nu_\tau + Fe & \rightarrow \tau + Fe^* \quad \text{(LSD)} \\
\bar{\nu}_e + Fe & \rightarrow \bar{\nu}_e + Fe^* \quad \text{(LSD)}
\end{align*}
\]

Here and in what follows \( l = e, \mu, \tau; \ x = e, \mu, \tau, \) and superscript "\(^*\)" denotes the exited states of the nuclei, which immediately decay to ground ones emitting nucleons and gammas. The role of reactions which involve exited states of nuclei for neutrino detection is discussed in detail elsewhere.\(^{7,7}\)

In 1987 only reaction \((??)\) was regarded to allow supernova neutrino detection. However, effective electron antineutrino detection area for KII exceeded those for LSD substantially. Numerous studies\(^{7,9,10,11,12}\) indicated that LSD signal could hardly be attributed to \((7-12)\) MeV electron antineutrino flux. We focus our attention on the other possibility, which was elaborated\(^{7,7}\) only in 2004: the idea was that neutrinos (not antineutrinos) of sufficiently high \((30-50\) MeV\) energy produced the signal in LSD through the reactions on iron and carbon nuclei. Effective LSD and KII areas for \( \nu_\tau \) and \( \bar{\nu}_e \) detection are given in Table 2. Cross sections for reactions \((??)\) are tabulated elsewhere.\(^{7,7}\)

One could infer from Table 2 that no contradiction between LSD and KII event numbers would occur if supernova neutrino flux was composed of electron and non-electron neutrinos of appropriate energies. However the question is weather such flux content is possible in principle. To answer this question one should consider supernova neutrino emission and flavour conversion.

Neutrinos and antineutrinos of all three flavours can be created during the collapse of the iron core in the centre of the star. All reactions in which they are created conserve lepton flavour. Moreover, muon neutrinos are produced in the same reactions as tau-neutrinos. Therefore in any collapse model neutrino fluxes satisfy the following conditions:

\[
F_{\nu_e} - F_{\bar{\nu}_e} \leq \frac{N_{\nu_e}^{Core}}{4\pi L^2} \leq \frac{N_{\nu_{\mu \tau}}^{Core}}{2M_{\odot}} \cdot 10^{10} \text{ cm}^{-2}
\]
Here $F_{\nu_{e},\bar{\nu}_{e}}^{0}$ is the time- and energy-integrated flux of an (anti-)neutrinos. Upper index "0" denotes that it is an original flux, i.e., such a flux which would reach the earth if there were no flavour conversion in the matter of the star. $M_{\text{Core}} = (1.4-2.2) M_{\odot}$ is the mass of the iron core, $N_{\text{eCore}}^{0}$ is the number of electrons in the core and $L = 52$ kpc is the distance between the supernova and the earth.

Fluxes at the earth $F_{\nu_{\mu},\bar{\nu}_{\mu}}$ are linear combinations of original fluxes:

$$
F_{\nu_{e}} = p F_{\nu_{e}}^{0} + (1-p) F_{\bar{\nu}_{e}}^{0} \quad F_{\bar{\nu}_{e}} = (1-p) F_{\nu_{e}}^{0} + (1+p) F_{\bar{\nu}_{e}}^{0}
$$

$$
F_{\nu_{\mu}} = p F_{\nu_{\mu}}^{0} + (1-p) F_{\bar{\nu}_{\mu}}^{0} \quad F_{\bar{\nu}_{\mu}} = (1-p) F_{\nu_{\mu}}^{0} + (1+p) F_{\bar{\nu}_{\mu}}^{0}
$$

(4)

Coefficients $p$, $\bar{p}$ (see Table 3) depend on the unknown neutrino mass hierarchy and neutrino mixing angle $\theta_{13}$.

Upper bounds on original fluxes $F_{x}^{0}$ and $F_{\nu_{x}}^{0}$ follow immediately from eq. (??) and Table 3:

$$
F_{x}^{0} \leq \frac{1}{1-p} F_{\nu_{x}} \leq 3.6 F_{\nu_{x}} \quad F_{\nu_{x}}^{0} \leq \frac{1}{p} F_{\nu_{x}}
$$

(5)

1 Results

We numerically investigated $P(F_{\nu_{x}}^{0}, F_{\bar{\nu}_{x}}^{0}, F_{x}^{0})$, the probability that fluxes $F_{\nu_{x}}^{0}$, $F_{\bar{\nu}_{x}}^{0}$ and $F_{x}^{0}$ after the flavour conversion according to eq. (??) produced not less than 5 events in LSD and not more than 2 events (with energies less than 12-14 MeV) in KII (the details are given elsewhere). Possible effects due to shock wave propagation were taken into account. The results proved to be stable under the reasonable variation of the input cross sections and detector efficiencies. They lead to the following conclusions concerning the first stage of the two-stage SN1987A explosion models.

(1) In the case of small neutrino 1-3 mixing angle, $\theta_{13} < 0.003$, two-stage SN1987A explosion models are disfavoured by the data, independently of the neutrino mass hierarchy.

(2) Non-electron neutrino and antineutrino production had to be severely suppressed during the first stage of the collapse, independently of the neutrino mass hierarchy and mixing angle $\theta_{13} :$

$$
F_{x}^{0} \lesssim 10^{8} \text{cm}^{-2}.
$$

(6)

This means that at the first stage of the collapse there was no thermal equilibrium, even rough.

(3) In the case of normal mass hierarchy and large 1-3 mixing angle, $\theta_{13} > 0.03$, in order to
explain the data one should imply emission of very energetic ($E \gtrsim 60$ MeV) electron neutrinos, $\nu_e$, at the first stage of the explosion; at the same time the suppression of $\bar{\nu}_e$ production should be assumed:

$$F_{\nu_e}^0 \simeq (0.3 - 0.5) \cdot 10^{10} \text{cm}^{-2}, \quad F_{\bar{\nu}_e}^0 \lesssim 10^8 \text{cm}^{-2}. \quad (7)$$

In addition, large values of the collapsing core mass, $M_{\text{Core}} \gtrsim 2M_\odot$, were necessary. A powerful shock wave could further complicate the agreement of the data with the theory.

(4) In the case of inverted mass hierarchy and large 1-3 mixing angle, $\theta_{13} > 0.03$, the data can be explained by the moderate energy ($50 \text{ MeV} \lesssim E \lesssim 45 \text{ MeV}$) electron neutrino and antineutrino emission at the first stage of the explosion with fluxes of order of $10^{10} \text{cm}^{-2}$. A powerful shock wave could worsen the agreement of the data with the theory.

Acknowledgments

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The role of lepton flavours in Thermal Leptogenesis

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Thermal leptogenesis, that can be viewed as a consequence of the see-saw model, is a very natural mechanism to explain the matter asymmetry of the Universe. Recently, lepton flavours have been included in the Boltzmann equations, modifying significantly the evaluation of the efficiency of leptogenesis to explain the observed baryon asymmetry.

1 Introduction

Among the weakness of the Standard Model (SM), two at least could be explained by the fact that neutrinos are massives. Those masses, deduced from neutrino oscillations experiment, as well as the mixings between the different generation of neutrinos, are well described within the see-saw mechanism. For that purpose, at least two right handed neutrinos of Majorana type that are singlet under the SM gauge group are introduced. Their mass scale, constrained by the seesaw model, can be much higher than the electroweak scale, in the case of a natural theory, that we consider in this proceeding. Leptogenesis, in the thermal scenario, is the creation of a net lepton number in the very early Universe from the decay of heavy right-handed neutrinos, that is (partially) converted into a baryon asymmetry through sphalerons processes.

1.1 The Baryon Asymmetry of the Universe

The baryon asymmetry, that is the difference between the number density of baryon and antibaryon normalised to the number density of photons, is deduced from cosmological observation to be:

$$ n_B = \frac{n_b - n_{\bar{b}}}{n_{\gamma}} = (6.1 \pm 0.2) \times 10^{-10}. $$

To explain this asymmetry in a inflationary universe, three conditions given by Sakharov are necessary: baryon number violation, C and CP violation, and departure from thermal equilib-
rium. These conditions are satisfied in the Standard Model, but hardly enough to explain the amount of baryon asymmetry. One needs to go beyond the SM.

1.2 The Seesaw mechanism

In the seesaw model\(^3\), neutrinos of Majorana type are added to the SM particle content to explain the smallness of the observed masses. The extended Lagrangian, in the mass basis of the right-handed neutrino \( N \) and of the charged lepton \( \ell \), reads:

\[
\mathcal{L} = \mathcal{L}_{\text{SM}} - \frac{1}{2} N_i M_{ij} N_j - h^{\alpha \beta}_i \ell^\alpha \bar{\ell}^\beta \phi .
\]  

(2)

\( M_{ij} \) is a mass matrix, diagonal in the basis we choose, and \( h_i^\alpha \) is a complex Yukawa coupling. The diagonalisation of the neutrino mass matrix gives two mass eigenstate (per generation): one with a mass \( \simeq M \), and the other with the mass \( m_\nu \simeq v^2 h_i M^{-1} h_i \), where \( v \) is the vacuum expectation value of the Higgs field \( \phi \). For light neutrino mass of about 1 eV, and natural couplings \( h_i \), \( M \) should be around \( 10^9 \text{GeV} \) or more.

2 Thermal Leptogenesis

In the thermal leptogenesis scenario\(^4\), the right handed neutrino \( N \) are produced in the thermal bath by scattering processes that occur at a temperature \( T \simeq M \). These \( N \) decay into lepton (plus higgs) and anti-lepton (plus higgs*), and if the decay violate CP a net lepton number is produced. Depending on the strength of inverse reaction, that is, depending on if inverse reaction are out-of-equilibrium or not, an asymmetry will survive, and will be partly converted into a baryon asymmetry by the \( B + \mathcal{L} \) - fast violating processes that are in equilibrium above the electroweak phase transition.

2.1 The standard picture: the one-flavour approximation

As in seesaw models leptogenesis qualitatively occurs, the first goal was to show that leptogenesis can quantitatively account for the observed matter-antimatter asymmetry. This has been done in the so called one-flavour approximation, where the lepton asymmetry is produced in one dominant flavour from the decay of the lightest right-handed neutrino, assuming a strong hierarchy: \( M_{N_1} \ll M_{N_2}, \cdots, M_{N_N} \). In this picture, it has been shown that leptogenesis can quantitatively works\(^5\). The baryon asymmetry is estimated to be:

\[
Y_B = \frac{n_B - n_\bar{B}}{s} \simeq a_{\text{sph}} \times Y^\text{eq}_{N_1} (T \gg M_{N_1}) \times \epsilon_{N_1} \times \eta ,
\]  

(3)

where \( a_{\text{sph}} \sim 1/3 \) is the fraction of lepton asymmetry converted into a baryonic one through the sphaleron interactions, \( Y^\text{eq}_{N_1} (T \gg M_{N_1}) \sim 4 \times 10^{-3} \) is the equilibrium number density of \( N_1 \) at the beginning of the leptogenesis era after the reheating period. The CP violation in the decay is parametrized by the CP asymmetry \( \epsilon_{N_1} \), which is defined as:

\[
\epsilon_{N_1} = \frac{\Gamma(N_1 \rightarrow H\ell) - \Gamma(N_1 \rightarrow H\bar{\ell})}{\Gamma(N_1 \rightarrow H\ell) + \Gamma(N_1 \rightarrow H\bar{\ell})}.
\]  

(4)

The last factor, \( \eta \), is the efficiency of the process, and highlights the competition between the production of a lepton asymmetry by decay and its destruction by inverse reaction (inverse decay and scattering processes). If the inverse reaction are fast compare to the Hubble expansion rate at the temperature \( T \simeq M_{N_1} \), then the lepton asymmetry will be strongly wash-out and not enough baryon asymmetry will be created.
2.2 The lepton flavour

The mass eigenstates of the particle contributing to the Boltzmann equations (BE) are determined by the reactions which are fast compared to the processes included in the BE. But at the temperature leptogenesis occurs, $T_{\text{lepto}} \simeq M_{N_1} \simeq 10^{12}$ GeV, as the fields acquire large thermal mass, the interaction involving charged Yukawa couplings develop thermal corrections:

$$\Gamma_{\ell_\alpha} \simeq 5 \times 10^{-3} \times h_{\ell_\alpha}^2 \times T.$$  

(5)

Depending on the Yukawa couplings $h_{\ell_\alpha}$, the interactions can be faster than the Hubble expansion rate at $T_{\text{lepto}}$, and have to be taken into account in the calculation of the proper mass eigenstates. More precisely, as $T_{\text{lepto}} \simeq M_{N_1}$, if $M_{N_1}$ is above about $10^{12}$ GeV, none of the interaction that bear the flavour information is in equilibrium, thus making indistinguishable the different flavour: the one-flavour approximation is valid. But if $M_{N_1}$ is below $10^{12}$ GeV the tau-Yukawa interactions are in equilibrium and two flavour are distinguishable: the flavour $\tau$ and an orthogonal flavour compose of $\mu$ and $e$. If $M_{N_1} \leq 10^9$ GeV the muon-Yukawas are in equilibrium too and the lepton asymmetry is projected onto a three flavour-space $\tau, \mu, e$.

3 Flavoured leptogenesis

In thermal leptogenesis the constraint that the reheating temperature should be above than $M_{N_1}$ in order not to wash-out the produced baryon asymmetry, but $T_{rh}$ should also not being too big, in order not to overdosure problem. Therefore, the lowest $M_{N_1}$ is the preferred choice. On the other hand, a lower bound on $M_{N_1}$ has been derived in the one flavour approximation \cite{7}, $M_{N_1} \geq 10^9$ GeV (in the case where $N_1$ is produced by thermal scatterings). As explained before, the flavour content should be taken into account: flavour matters \cite{8}. We thus have to define a CP asymmetry for each (distinguishable) flavour,

$$\epsilon_{N_1, \ell_\alpha} = \frac{\Gamma(N_1 \rightarrow H\ell_\alpha) - \Gamma(\bar{N}_1 \rightarrow \bar{H}\bar{\ell}_\alpha)}{\Gamma(N_1 \rightarrow H\ell_\alpha) + \Gamma(\bar{N}_1 \rightarrow \bar{H}\bar{\ell}_\alpha)},$$

(6)

as well as individual efficiencies $\eta_\ell$, so that the baryon asymmetry, when flavours are accounted for, reads:

$$Y_B \simeq a_{\text{sph}} \times Y_{N_1}^{\text{eq}}(T \gg M_{N_1}) \times \sum_\alpha \epsilon_{N_1, \ell_\alpha} \times \eta_\ell.$$  

(7)

Recall that in the one-flavour approximation we have $\sum_\alpha \epsilon_{N_1, \ell_\alpha} = \eta_\ell$, where $\eta_\ell$ is the efficiency factor for the total lepton asymmetry. This comes from the fact that in the Boltzmann equations for the number densities, each (distinguishable) flavour is washed-out with a strength $\propto \bar{m}_\alpha/m^*$, where $\bar{m}_\alpha$ is the rescaled partial decay width $\Gamma(N_1 \rightarrow H\ell_\alpha)$ and $m^* \simeq 10^{-3}$eV is the "equilibrium neutrino mass", the rescaled Hubble expansion rate at $T_{\text{lepto}}$. In the one-flavour case, the total lepton asymmetry is washed-out with a strength $\propto \sum_\alpha \bar{m}_\alpha/m^*$, then possible flavour misalignment can enhanced the amount of lepton asymmetry. Indeed, the efficiency is maximum for $\bar{m} \simeq m^*$, but the mass inferred from neutrino oscillations, $m_{\text{atm}} \simeq 5 \times 10^{-2}$eV and $m_{\text{sol}} \simeq 9 \times 10^{-3}$eV are both above $m^*$, and thus the region $\bar{m} \geq m^*$ is more attractive, even if less efficient for leptogenesis. Including flavour, one can have $\sum_\alpha \bar{m}_\alpha \geq m^*$, even if one of the flavour is weakly washed-out $\bar{m}_\alpha \simeq m^*$. The efficiency for this flavour will be (close to) maximum, and the flavour will dominate the lepton asymmetry (unless its CP asymmetry $\epsilon_{N_1, \ell_\alpha}$ is too small...), allowing for a sufficient amount of baryon asymmetry, even if the total wash-out is strong $\sum_\alpha \bar{m}_\alpha \gg m^*$.

Another feature of the inclusion of lepton flavour concern the upper bound on the light neutrino mass scale. In the one-flavour case, an upper bound of 0.15 eV was derived \cite{9} on the light neutrino mass scale, from the requirement that the total wash-out should not be too strong. But as we have seen, flavour misalignment allows successful leptogenesis even for a
strong total wash-out if one flavour is in weak or mild wash-out $\bar{m}_0 \simeq m^*$, and therefore no upper bound on $m_\nu$ holds from leptogenesis: the cosmological bound is saturated.

4 Conclusion

In seesaw models leptogenesis qualitatively occurs: a right handed neutrino produce a lepton asymmetry via out-of-equilibrium CP violating decays. It has been shown that it quantitatively explain the observed amount of baryon asymmetry in the one-flavour approximation. Including lepton flavours, the computation of the baryon asymmetry is modified, and some constraint are relaxed, from possible mis-alignment of the flavours. For the good range of temperature, flavoured treatment of leptogenesis is more accurate.

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References

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