2007 QCD and High Energy Hadronic Interactions
2007 QCD and High Energy Hadronic Interactions

edited by

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Bolek Pietrzyk
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The QCD Session of the XLIII\textsuperscript{nd} Rencontres de Moriond

was organized by

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with the active collaboration of:

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B. Pietrzyk,
C. I. Tan,
J. Trân Thanh Vân,
U. Wiedemann.
The XLIInd 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 XLIInd 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:


E. Augé, E. Berger, S. Bethke, A. Capella, A. Czarnecki, D. Denegri, Y. Dokshitzer, N. Glover, B. Klima, L. Kluberg, M. Krawczyk, L. McLerran, B. Pietrzyk, Chung-I Tan and U. Wiedemann for the «QCD and High Energy Hadronic Interactions» session,

Kevin Baines, Ludwik Celnikier, Pierre Drossart, David Grinspoon, Håkan Svedhem, Frederik Taylor, Dmitri Titov, Olivierie 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.

J. Trần Thanh Vân
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I - Mini session on fits

Chairpersons: B. Pietrzyk
Introduction to Frequentist and Bayesian Fitting

F. James

CERN, Geneva, Switzerland

A brief introduction to Bayesian and frequentist statistical analyses, emphasising the differences between the two.

Frequentist Probability

Frequentist statistics is based on frequentist probability, which is defined as a limiting frequency, over a set of (hypothetical) repetitions of the same experiment.

\[ P(A) = \lim_{N \to \infty} \frac{N(A)}{N} \]

where \( A \) occurs \( N(A) \) times in \( N \) trials.

Frequentist probability is used in most scientific work, because it is objective. It can (in principle) be determined to any desired accuracy and is the same for all observers.

It is the probability of Quantum Mechanics.

It can only be applied to repeatable phenomena.

Frequentist Fitting

An important principle in frequentist statistics is that you should always be able to calculate how often you are wrong:

- How often the quoted error bar does not include the true value.
- How many unwanted events will pass the cuts and contaminate the sample.
- How many good events will be rejected by the cuts and reduce the signal.
- How often the Chi-square Test will fail even if the hypothesis is correct.

The hard part of frequentist statistics is to find methods such that the probability of being wrong is independent of the value(s) of any unknown parameter(s).

For confidence intervals, J. Neyman (with the help of Gary Feldman and Bob Cousins) showed us how to do this.
Bayesian Probability

Bayesian Statistics is based on the Bayesian definition of Probability which is a degree of belief.

The operational definition most often used is de Finetti’s coherent bet. Your belief in $A$ is proportional to the amount you would bet on $A$ happening.

This allows you to define the probability of non-repeatable events.

In particular, the prior probability of $A$ is your belief in $A$ before you do any experiments to measure $A$.

Bayesian Fitting

All Bayesian methods require as input a prior probability. This is a kind of phase space: All experimental results are multiplied by the prior to obtain the posterior probability.

Concerning the prior, there is some good news and some bad news:

1. The bad news is:
   - Unlike physical phase space, there is no principle that tells you what prior probability to use. It is arbitrary.
   - No matter what your experiment tells you, it is always possible to find a prior that gives any result you want.

2. The good news is: The prior becomes less important as the amount of data increases. However, even this has to be taken cautiously because:
   - The limit in which the data dominate the prior is anyway the limit in which also frequentist methods give the same results.
   - When the hypothesis is multidimensional (several parameters estimated simultaneously), it cannot be shown that the data must dominate the prior.

So Why would any Physicist Use Bayesian Methods?


But some physicists are Bayesian. Why?

- Statisticians consider Bayesian methods valid, and in some sense even more coherent than frequentist methods.
- Bayesian has become “trendy”, “modern”.
- Bayesian is easier, more natural.
- The Karmen Problem.
The Karmen Problem before Feldman-Cousins

The Karmen Problem is the observation of more events (signal plus background) than are expected from background alone, making the apparent observed signal negative. Before the work of Feldman and Cousins, the standard frequentist upper limit in this case was given by the table below, which shows that it could be negative. This method would still have overall exact coverage, since the probability of such a result is very small.

### Frequentist 90% Upper Limits for Poisson with Background

<table>
<thead>
<tr>
<th>observed</th>
<th>0</th>
<th>1</th>
<th>2</th>
<th>3</th>
</tr>
</thead>
<tbody>
<tr>
<td>background = 0.0</td>
<td>2.30</td>
<td>3.89</td>
<td>5.32</td>
<td>6.68</td>
</tr>
<tr>
<td>0.5</td>
<td>1.80</td>
<td>3.39</td>
<td>4.82</td>
<td>6.18</td>
</tr>
<tr>
<td>1.0</td>
<td>1.30</td>
<td>2.89</td>
<td>4.32</td>
<td>5.58</td>
</tr>
<tr>
<td>2.0</td>
<td>0.30</td>
<td>1.89</td>
<td>3.32</td>
<td>4.68</td>
</tr>
<tr>
<td>3.0</td>
<td>-0.70</td>
<td>0.89</td>
<td>2.32</td>
<td>3.68</td>
</tr>
</tbody>
</table>

The naive Bayesian upper limit (with uniform prior) for the same case gave just what the physicist wanted:

### Bayesian 90% Upper Limits (Uniform Prior)

<table>
<thead>
<tr>
<th>observed</th>
<th>0</th>
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<tbody>
<tr>
<td>background = 0.0</td>
<td>2.30</td>
<td>3.89</td>
<td>5.32</td>
<td>6.68</td>
</tr>
<tr>
<td>0.5</td>
<td>2.30</td>
<td>3.50</td>
<td>4.83</td>
<td>6.17</td>
</tr>
<tr>
<td>1.0</td>
<td>2.30</td>
<td>3.26</td>
<td>4.44</td>
<td>5.71</td>
</tr>
<tr>
<td>2.0</td>
<td>2.30</td>
<td>3.00</td>
<td>3.87</td>
<td>4.92</td>
</tr>
<tr>
<td>3.0</td>
<td>2.30</td>
<td>2.83</td>
<td>3.52</td>
<td>4.37</td>
</tr>
</tbody>
</table>

The Karmen Problem after Feldman-Cousins

Feldman and Cousins discovered that we had not been applying frequentist statistics in the correct way, and it should be done using the maximum likelihood ratio ordering principle, which would assure a unified approach to upper limits and two-sided limits, while at the same time eliminating negative upper limits.

### Feldman-Cousins 90% Upper Limits for Poisson with Background

<table>
<thead>
<tr>
<th>observed</th>
<th>0</th>
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<th>3</th>
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<tbody>
<tr>
<td>background = 0.0</td>
<td>2.44</td>
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<tr>
<td>0.5</td>
<td>1.94</td>
<td>3.86</td>
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<tr>
<td>1.0</td>
<td>1.61</td>
<td>3.36</td>
<td>4.91</td>
<td></td>
</tr>
<tr>
<td>2.0</td>
<td>1.26</td>
<td>2.53</td>
<td>3.91</td>
<td>5.42</td>
</tr>
<tr>
<td>3.0</td>
<td>1.08</td>
<td>1.88</td>
<td>3.04</td>
<td>4.42</td>
</tr>
</tbody>
</table>

At the same time, people questioned the uniform Bayesian prior, and found the the Jeffreys prior had a better theoretical base and corresponded better to actual prior belief. Unfortunately, the Jeffreys prior gives:

### Bayesian 90% Upper Limits (Jeffreys Prior)

<table>
<thead>
<tr>
<th>observed</th>
<th>0</th>
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<tbody>
<tr>
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<td>3.89</td>
<td>5.32</td>
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<tr>
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<tr>
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<td>2.0</td>
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<td>0.00</td>
</tr>
<tr>
<td>3.0</td>
<td>0.00</td>
<td>0.00</td>
<td>0.00</td>
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</tr>
</tbody>
</table>
Therefore, the Karmen problem, or more generally the problem of getting reasonable upper and lower limits (with correct frequentist coverage) in delicate cases with very small numbers of events, is no longer a reason to adopt Bayesian statistics. On the contrary, recent advances show the exact frequentist methods to be closer to what the physicist wants.

What can go wrong with Bayesian analyses?

One of the most striking examples of malfunctioning of data analysis due to the use of Bayesian principles can be found in the paper “Evidence for neutrinoless double beta decay”, by Klapdor-Kleingrothaus et al, submitted to World Scientific on May 20, 2006. This research article finds strong statistical evidence for a peak where many physicists don’t see any peak. One of the criticisms of the analysis was that the authors did not do a standard goodness-of-fit test for the no-peak hypothesis to see whether the data required a peak. But the standard chi-square test is not allowed in the Bayesian framework because it violates the Likelihood Principle. The explanation given by the authors was as follows:

Criticism: There is no null hypothesis analysis demonstrating that the data require a peak.

Reply: The statement that there is a peak with probability K implies that there may be none. Since the results are probabilistic, it is not possible to demonstrate that the data do require a peak.

Note that the model used to analyze the data includes the possibility that the intensity of the line is zero. The error intervals given in KDHK have been chosen such as to exclude the value of zero. Therefore every result in KDHK can especially be read as: “With probability K the intensity zero is outside the error interval”. In this sense the null hypothesis is rejected with the probability on which the error interval is based.

The difficulty we have in understanding this reply makes us question whether Bayesian methods are really so natural and intuitively appealing.

Conclusions: Bayesian or frequentist?

1. The main problem in Bayesian methodology is the prior. Use Bayesian methods when you know the prior and have a good reason to use it. The only case I know where that is true is maximum entropy image processing.

2. Use Bayesian decision theory to make it clear what are the subjective criteria for your decision. [Example: where to look for new physics.]

3. For everything else, in particular objective data analysis, I don’t see any reason to use Bayesian methods. We now know how to handle all the situations (nuisance parameters, systematic errors) that used to cause problems in the frequentist methodology.

4. Very few people would believe a result that can only be obtained by a Bayesian analysis with an arbitrary prior.
BAYESIAN STATISTICAL METHODS IN PARTICLE PHYSICS

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Traditionally data analysis in High Energy Physics has relied on frequentist statistical methods, where probability is associated only with possible outcomes of repeatable observations. Often, however, we need to quantify uncertainties that are not easily described in this framework, such as model uncertainties or other systematic errors. Here we examine how Bayesian statistical methods can be used in such circumstances.

1 Introduction

The goal of a statistical analysis in particle physics is to make inferences about models or parameters using information from a finite data sample. The ‘statistical’ or sampling error is an inescapable element of uncertainty that characterizes any result, and it can be quantified using standard tools of frequentist statistics. Here a probability is associated with the possible outcomes of a repeatable observation. The main tools in this approach include statistical hypothesis tests and parameter estimation using, e.g., the method of maximum likelihood.

Although frequentist statistics remains the primary method of data analysis in HEP, it cannot answer a number of relevant questions, such as “What is the probability that my parameter is contained in such-and-such a fixed range?” or “What is the probability that this-or-that model is true?” In the frequentist framework, such probabilities are either 0 or 1, although it may not be known which. In Bayesian statistics, however, we extend the interpretation of probability to include degree of belief, also called ‘subjective probability’, and in this way an answer of 50% to the previous questions could be perfectly reasonable. Bayesian methods will thus be especially useful for treating model uncertainties and other types of systematic errors that are assumed not to fluctuate upon repetition of the measurement.

The basic idea behind Bayesian statistics is to interpret the probability of a hypothesis $H$ as a degree of belief that the hypothesis is true. We use Bayes’ theorem to calculate the conditional probability of $H$ given the measurement $x$. 
\[ p(H|x) = \frac{p(x|H)p(H)}{\int p(x|H)p(H) \, dH} \quad (1) \]

Here \( x \) denotes whatever collection of numbers characterizes the outcome of our measurement. The hypothesis \( H \) will often refer to a parameter value. The quantity \( p(H) \) is the prior probability for the hypothesis, i.e., it reflects one's degree of belief about \( H \) before having seen the outcome of the experiment \( x \). This multiplies the likelihood \( p(x|H) \), which is the probability for the data \( x \) under the assumption of the hypothesis \( H \). The denominator in (1) is an integral (or sum) over all possible hypotheses, and serves to normalize the probability \( p(H|x) \) to unity. The combination of ingredients on the right-hand side of (1) gives \( p(H|x) \), the posterior probability for \( H \) given the outcome \( x \). Thus Bayes' theorem tells us how to update our degree of belief about \( H \) in the light of the data \( x \).

2 Example of a Bayesian fit

We can compare the frequentist and Bayesian approaches using a simple example. Suppose we have \( n \) measured data values \( y_i, i = 1, \ldots, n \), corresponding to \( n \) values of a (known) control variable \( x \). Suppose further that the \( y_i \) are modelled as independent Gaussian random variables with known standard deviations \( \sigma \) and expectation (i.e., true) values \( \mu(x) \). Our goal is to fit the function \( \mu(x) \) to the data.

Suppose now the function \( \mu(x) \) is a straight line, \( \mu(x; \theta_0, \theta_1) = \theta_0 + \theta_1 x \), but we are only interested in the parameter \( \theta_0 \). The slope \( \theta_1 \) is an example of a nuisance parameter; it is needed in the fit, but its value is of no interest at the end of the analysis. At least formally, systematic uncertainties can be viewed as an uncertainty in some number of nuisance parameters.

In frequentist statistics, we estimate \( \theta_0 \) and \( \theta_1 \) by finding the values that maximize the likelihood function,

\[ L(\theta_0, \theta_1) = \prod_{i=1}^{n} \frac{1}{\sqrt{2\pi}\sigma_i} \exp \left[ -\frac{1}{2} \frac{(y_i - \mu(x_i; \theta_0, \theta_1))^2}{\sigma_i^2} \right] \quad (2) \]

Equivalently in this case we can minimize the quantity \( \chi^2(\theta_0, \theta_1) = -2 \ln L(\theta_0, \theta_1) \). The estimators \( \hat{\theta}_0 \) and \( \hat{\theta}_1 \) are characterized by standard deviations \( \sigma_{\theta_0} \) and \( \sigma_{\theta_1} \), which can be obtained from the tangents to the contour where the log-likelihood function decreases by 1/2 from its maximum value, or equivalently where the \( \chi^2 \) increases by one unit from its minimum (see, e.g., [1]). Because in general the estimators of the parameters will be correlated, the standard deviation of the parameters of interest winds up being inflated as a result of the statistical uncertainty in the nuisance parameters.

It could be, however, that we have additional information about the nuisance parameter \( \theta_1 \) through a subsidiary measurement \( t_1 \), which we might model as a Gaussian random variable with a standard deviation \( \sigma_{t_1} \). Including this information in the likelihood function, or equivalently adding a term \( (t_1 - \theta_1)^2 / \sigma_{t_1}^2 \) to the \( \chi^2 \), provides not only a better determination of the nuisance parameter \( \theta_1 \), but also it improves the measurement of the parameter of interest \( \theta_0 \). This is illustrated in Fig. 1(a), which shows the contours of constant \( \chi^2 \) from which one determines the standard deviations both with and without the additional measurement \( t_1 \).

In order to solve the same problem using the Bayesian approach, we need first to write down a joint prior probability density \( p(\theta_0, \theta_1) \) for the two parameters in the problem. For this example let us suppose that this prior factorizes into a product of terms \( \pi(\theta_0, \theta_1) = \pi_0(\theta_0)\pi_1(\theta_1) \). For the parameter \( \theta_0 \) where we have the subsidiary measurement \( t_1 \), take \( \pi_1(\theta_1) \) to be a Gaussian in \( \theta_1 \) centered about our measured value \( t_1 \) and having a standard deviation \( \sigma_{t_1} \). For \( \theta_0 \), suppose we have essentially no prior information and so we choose its prior to be much broader than the
likelihood: often it is convenient to set this prior equal to a constant $\pi_0$ (the technical difficulty with non-normalizability is not important in this example). We then put these ingredients into Bayes’ theorem to find the joint posterior pdf for $\theta_0$ and $\theta_1$ given the measurements $\vec{y}$.

$$p(\theta_0, \theta_1 | \vec{y}) \propto \prod_{i=1}^{n} \frac{1}{\sqrt{2\pi \sigma_i}} e^{-\frac{(y_i - \mu(y_i, \theta_0, \theta_1))^2}{2\sigma_i^2}} \pi_0 \frac{1}{\sqrt{2\pi \sigma_1}} e^{-\frac{(\theta_1 - \xi_1)^2}{2\sigma_1^2}}.$$ \hspace{1cm} (3)

Here we write Bayes’ theorem as a proportionality, since the denominator in (1) is determined by the requirement that the probability density function (pdf) be normalized to unity. Since we are only interested in the parameter $\theta_0$, we marginalize (integrate) over the nuisance parameter $\theta_1$.

$$p(\theta_0 | \vec{y}) = \int p(\theta_0, \theta_1 | \vec{y}) d\theta_1$$ \hspace{1cm} (4)

In this simple example one can carry out the integral in closed form. We find in fact an answer which appears to coincide exactly with what we obtained using the frequentist approach, namely, a Gaussian centred about the point where the likelihood (including the additional measurement $t_1$) is maximized, and having a covariance matrix equal to that of the estimators $\theta_0$ and $\theta_1$ from the frequentist method of maximum likelihood.

It appears therefore that the Bayesian method has so far bought us very little, other than the fact that we may now interpret the pdf $p(\theta_0 | \vec{y})$ as a function representing our degree of belief about where $\theta_0$ lies. The real advantage of the Bayesian method, however, becomes apparent when we wish to include prior information that does not stem directly from measurements but which could come, say, from arguments of symmetry or theoretical prejudices.

Suppose, for example, that the nuisance parameter $\theta_1$ represents a quantity such as a coefficient in a perturbation series that could in principle be computed, but the calculation has not yet been carried out. Our theorist friend who has experience in these things might nevertheless have a general feeling about how large the value may be; it might be highly certain that its value is positive and to ‘probably’ be of the order of 0.1 or thereabouts. Under pressure, the theorist sketches a prior that looks roughly like an exponential curve,

$$\pi_1(\theta_1) = \frac{1}{\tau} e^{-\theta_1/\tau},$$ \hspace{1cm} (5)

for $\theta_1 > 0$ and zero otherwise, with the parameter $\tau = 0.1$. In order to extract this information from the theorist it may prove useful to recall the ‘if-then’ nature of Bayes’ theorem: if one assumes the prior (5) then Bayes’ theorem tells us how these beliefs should be updated in the
light of new data. In order for a Bayesian analysis to be of value to a broader community of scientists, who may have different prior beliefs, the analyst should show how the posterior probabilities change upon a reasonable variation of the prior. This is indicated in Fig. 1(b), which shows the marginalized posterior pdf for the parameter of interested \( p(\theta | y) \) obtained using three different values of the parameter \( r \). In this example, the data points are such that a smaller value for the slope \( \theta_1 \) implies a larger value for the intercept \( \theta_0 \) (i.e., the \( g_i \) increase with increasing \( r_i \)). In a similar way one could try different functional forms for the prior.

It is worth noting that the integral (4) that one carries out to marginalize over the nuisance parameters can rarely be done in closed form. In cases with large numbers of parameters it can not even be computed using usual Monte Carlo methods (e.g., acceptance-rejection). The field of Bayesian computation has made great progress in recent years exploiting techniques such as Markov Chain Monte Carlo (MCMC); see e.g., 2. This method generates a correlated sequence of points in the parameter space, so the effective sampling error will decrease more slowly as a function of the number of values \( n \) than the usual \( 1/\sqrt{n} \). Nevertheless, the MCMC method allows one to find marginal pdfs for problems that would otherwise remain intractable with traditional Monte Carlo methods.

3 Conclusions

Sometimes we have prior information about parameters that stems directly from measurements for which we have a reliable probability model. This can be incorporated directly into the likelihood function and frequentist methods applied to estimate parameters or test hypotheses. At other times, however, we have information about parameters that is of a more subjective nature, e.g., when we have a certain degree of belief about where the true value of the parameter lies. In such cases this can be incorporated into our analysis using Bayesian methods.

Computational tools such as Markov Chain Monte Carlo allow one to marginalize joint pdfs over the nuisance parameters in order to obtain the pdf for the parameters of interest. Such techniques will prove valuable in HEP for quantifying, e.g., uncertainties that arise from an incomplete calculation of a theoretical prediction or an imperfect modelling of an experimental apparatus; examples related to uncertainties in parton densities are described in 3.

The reluctance to include subjectivity into a scientific analysis can be overcome by recognizing the ‘if-then’ nature of a Bayesian analysis. That is, one does not in general insist on a particular set of prior degrees of belief, but rather demonstrates how a range of reasonable prior beliefs evolve in the light of the data.

Acknowledgements

The author would like to convey his thanks to the organizers for a very enjoyable and fruitful meeting, and to Fred James and François Le Diberder for interesting and productive discussions about Bayesian and Frequentist statistics.

References

2. Markov Chain Monte Carlo methods are discussed in numerous texts, e.g., Phil Gregory, Bayesian Logical Data Analysis for the Physical Sciences, CUP, 2005.
II - Heavy flavour session

Chairpersons: B. Pietrzyk and J. E. Augustin
Rencontres de Moriond 2007
INTRODUCTION TO HEAVY FLAVOUR PHYSICS
for MoriondQCD 2007
(…more on CP violation and CKM physics)

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Laboratoire de l’Accélérateur Linéaire
IN2P3-CNRS et Université Paris Sud Sud 11,
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This paper presents an introduction to the heavy flavour physics underlying some interesting results which will be presented at this conference. Special emphasis is put on measurements related to CP violation and the determination of the Unitarity Triangle parameters.

Keywords: Flavour Physics, CP Violation, CKM matrix, Unitarity Triangle, New Physics

1 Introduction

Accurate studies of the production and decays of beauty and charm hadrons are exploiting a unique laboratory for testing the Standard Model in the fermion sector, for studying QCD in the non-perturbative regime and for searching for New Physics (NP) through virtual processes. Furthermore the main objective of the elementary particle physics today is to search for evidence of physics beyond the Standard Model. The production of new particles is not the only way to search for. New particles can manifest themselves through virtual effects in decays of Standard Model particles such as $B$ and $D$ mesons and $\tau$ leptons. The effects due to the presence of NP are proportional to the ratio coupling/scale. If the scale is fixed (particles are discovered) precise measurements in the flavour sector allow to measure the couplings and so to have a knowledge on the flavour structure of the NP. On the other hand precise measurements could allow us to probe New Physics energy scales inaccessible at present and next-generation colliders. The “mot d’ordre” for this approach is precision.

In the last few years tremendous improvements have been achieved with the result of a precise determination of the CKM parameters and rare decays from B-factories together with new results
from the Tevatron on B, physics. Many additional measurements on B and D mesons properties (masses, branching fractions, lifetimes...) are necessary to constrain the heavy Quark theories (Operator Product Expansion (OPE) / Heavy Quark Effective Theory (HQET) / Lattice QCD (LQCD)). The control of these theories is crucial for constraining the SM and looking for NP effects.

In the first part of this paper we show the new experimental results which are further constraining the Standard Model in the fermion sector, with emphasis on those which are the most sensitive to NP effects. We finally discuss the phenomenological impact of these measurements and we conclude on showing how to go on to try to use flavour physics as a very effective probe for NP.

2 Experimental results, mainly novelties as at Winter 2007.

Many new interesting results are presented to this conference, related to Unitarity Triangle (UT) angles, semi-leptonic and rare B decays and the recent evidence of the charm mixing. Few items are selected. Many others will be presented to this conference.

$\beta$ angle. The mixing induced CP asymmetry, $a_{J/\psi K_{s,L}}$, in $B^0_d \to J/\psi K_{s,L}$ decays allows to determine the angle $\beta$, with small theoretical uncertainties. The most recent measurements give a world average $^{1,2} \sin(2\beta) = 0.675 \pm 0.026$. One of the two solutions : $\beta = (21.2 \pm 1.0)^o$ is consistent with the Standard Model. The determination of $\sin 2\beta$ gives so far the best determination of $\rho$ and $\eta$ parameters. New results on $\beta$ are coming from the $B^0 \to D^{(*)} h^0$ decay mode. The determination of $\beta$ from these modes is not yet competitive with the previous one. Nevertheless together with $B \to J/\psi K \pi$ and $B \to D^* D^* K$ modes it helps in strongly disfavouring the solution of $\beta$ corresponding to $(68.1 \pm 1.0)^o$.

$\alpha$ angle. The angle $\alpha$ can be obtained using the time-dependent analyses of $B^0 \to \pi^+ \pi^-$. $B^0 \to \rho^+ \rho^-$ and $B^0 \to (\rho \pi)^0$. In the absence of contributions from penguin diagrams, these decays give a measurement of $\sin 2\alpha$. Several strategies have been proposed to control this so-called “penguin pollution”, which are mainly based on the use of an isospin analysis. Measurements of branching fractions and CP asymmetries have been made in both the $B \to \pi \pi$ and $B \to \mu \mu$ systems.

To better constrain the penguin contributions the measurement of the branching fractions for the neutral modes is crucial. For this a very important result has been obtained from Babar measuring $^3$: $Br(B^0 \to \rho^0 \rho^0) = (1.07 \pm 0.33 \pm 0.19) \times 10^{-6}$. It is a 3.5$\sigma$ evidence, and implies that the contributions from penguins in the $\mu \mu$ system are sizeable. Combining all the most recent results (including new results on $B \to \rho \pi$) we obtain $\alpha = [81, 111] \ U[159, 171]$ at 95% C.L. Selecting the SM solution: $\alpha = (93 \pm 8)^o$.

$\gamma$ angle. Various methods related to $B \to DK$ decays have been proposed to determine the UT angle $\gamma$. Using the fact that a charged $B$ can decay into a $D^0(\bar{D}^0)K$ final state via a $V_{ub} (V_{cb})$ mediated process. CP violation occurs if the $D^0$ and the $\bar{D}^0$ decay to the same final state. These processes are thus sensitive to the phase difference $\gamma$ between $V_{ub}$ and $V_{cb}$. Combining all the available measurements we get $\gamma = (82 \pm 20)^o$ with a $\pi$ ambiguity. The most precise results are coming from the Dalitz methods which consists in studying the interference between the $b \to u$ and the $b \to c$ transitions using the Dalitz plot of D mesons reconstructed into three-body final states. So far the only $D^0$ decay used was $D^0 \to K_s \pi^- \pi^-$. It will be important in future to control and reduce the error coming from the modelling of the D Dalitz decay. For this reason it is important to reconstruct other 3 bodies D decays. An important result from BABAR using...
the decay mode $D^0 \rightarrow \pi^- \pi^0 \pi^0$ gives $\gamma = (25 \pm 18)^\circ$.

$|V_{ub}|$. In a recent paper\textsuperscript{8} it was observed that the determination of $|V_{ub}|$, using inclusive methods, was disfavoured by all other constraints at the $2.5\sigma$ level. This can come either from the fact that the central value of $|V_{ub}|$ from inclusive decays is too large, or from the smallness of the estimated error, or both. On the other hands $|V_{ub}|$ from exclusive decays has still large uncertainties. Moreover the problem has been recently worsened by the decrease of the value of $\sin(2\beta)$ determined by the direct measurements. At this conference new results are presented on both determinations. An interesting approach consists on using the universality of the shape function to link the spectrum ($E_x$) of $B \rightarrow X_s \gamma$ decays with the one ($E_{lep}$, $M_X$) of $B \rightarrow X_u \ell \nu$ decays. Using the Babar data, it has been also shown that the value extracted for $|V_{ub}|$ is stable irrespective of the cuts on these observables and of the use of different theoretical approaches\textsuperscript{9}. In general the most recent analyses confirm high value for $|V_{ub}|$ (for example $|V_{ub}| = (1.40 \pm 0.30(stat) \pm 0.17(syst)) \times 10^{-3}$ as quoted in\textsuperscript{8}). This approach is interesting because part of the theoretical error is now absorbed in the statistical one.

Novelties also appeared on the exclusive analyses such as $(B \rightarrow \pi \ell \nu)$. The traditional solution to improve the S/B ratio was to cut on the quality of $p_{miss}$, to have a good reconstruction for the neutrino. With the increase of the available statistics new analyses can be also performed with much worse S/B ratio (of about 0.1), keeping severe criteria on the identification of $\pi - \ell$ pair. at relaxing the criteria on $p_{miss}$ and adjusting the cuts as a function of $q^2$. The results of that is an increase of more than a factor four in statistics. Babar has performed such analysis in 12 bins of $q^2$, obtaining $|V_{ub}| = (3.7 \pm 0.2 \pm 0.2^{+0.6}_{-0.1}(theo)) \times 10^{-3}$, which is the best single world determination of $|V_{ub}|$. The HFAG average\textsuperscript{1} is $|V_{ub}| = (3.55 \pm 0.22^{+0.6}_{-0.1}(theo)) \times 10^{-3}$. To improve this determination more precise calculations of the form factors are needed.

$\beta$ angle “with penguins”.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure1.png}
\caption{Left plot: HFAG compilation of measurements of $\sin(2\beta_{\text{eff}})$ in decays dominated by $b \rightarrow s$ penguin amplitudes. Right plot: Signal for $B^0 \rightarrow \rho^0 \pi^-$.}
\end{figure}

Very interesting channels to search for New Physics effects in mixing-induced $CP$ viola-
tion are those dominated by the $b \rightarrow s$ penguin transition. In the SM the measurements of $\sin 2\beta$ should be approximately the same as the one measured with charmonium B decays up to hadronic corrections of $\sim 0.02/0.05$ on $\sin(2\beta)$ (see, for example). New Physics particles in the loops could produce significant deviation from Standard Model predictions. Quite few new results are presented at the winter conferences, with many channels now precisely measured: $B^0 \rightarrow \phi K^0$, and $B^0 \rightarrow K^0 K^0 K^0$. For $B^0 \rightarrow \pi^+ K^0$, each individual channel gives a result on $\sin 2\beta$ smaller than the SM one and a naive average gives: $\sin 2\beta_{eff} = 0.53 \pm 0.05$ at about 2.5 $\sigma$ from the SM value. The results are shown in Fig. 1.

Radiative decays. The radiative FCNC decays $b \rightarrow s \gamma$ and $b \rightarrow d \gamma$ are very sensitive probes of New Physics. The ratio of rates of $b \rightarrow d \gamma$ and $b \rightarrow s \gamma$ decays can be used to determine $|V_{td}/V_{ts}|$. The theoretically cleanest case is for the neutral modes, because weak annihilation contributes more significantly in the charged mode.

The current experimental world averages are $\mathcal{B}(B^0 \rightarrow \rho^0 \gamma) = (0.91 \pm 0.19) \times 10^{-6}$ and the signal for $B^0 \rightarrow \rho^0 \gamma$ is shown in Fig. 1. From these results (using also the measured value of $\mathcal{B}(B^0 \rightarrow K^0 \gamma)$) we can extract $|V_{td}/V_{ts}| = 0.202 \pm 0.017 (\exp.) \pm 0.015 (\text{theor.})$. The theoretical error takes into account the ratio of $B \rightarrow V \gamma$ form factors and the remaining non-perturbative contributions.

Leptonic decays: $B \rightarrow \tau \nu$. Leptonic decay processes are described by annihilation diagrams and the rates of leptonic decays of the $B^+$ meson are proportional to $f_B^2 |V_{tb}|^2$, where $f_B$ is the pseudoscalar constant. This channel is interesting because it is sensitive to New Physics and in particular to charged Higgs exchange in a scenario with large tan $\beta$.

Signals for $B^+ \rightarrow \tau^+ \nu_\tau$ have been obtained, getting $B^+ \rightarrow \tau^+ \nu_\tau = (1.31 \pm 0.48) \times 10^{-4}$, which compares with the value expected in SM of $B^+ \rightarrow \tau^+ \nu_\tau = (0.85 \pm 0.13) \times 10^{-4}$. The experimental result has to be improved to become a significant test of new physics. Assuming the SM, the pseudoscalar constant can be extracted from this measurement to be $f_B = (237 \pm 37)$ MeV, which compares with $f_B = (189 \pm 27)$ MeV from lattice QCD calculations.

Evidence of $D$ mixing. The highlight of the conference is the new result on charm mixing, BABAR finds evidence for oscillations in $D^0 \rightarrow K^+ \pi^-$ with $3.9\sigma$ significance by studying the proper-time distribution for Wrong Sign (WS) data ($D^0 \rightarrow K^+ \pi^-$). The result is shown in Fig. 2.

Belle sees a $3.2\sigma$ effect in $D^0 \rightarrow K^+ K^-$, with results using $D^0 \rightarrow K^0 \pi^+ \pi^-$ supporting the claim. The world average of the mixing result gives:

$$x_D = (8.5_{-3.3}^{+4.2}) \times 10^{-3} \quad \text{and} \quad y_D = (7.1_{-2.2}^{+2.0}) \times 10^{-3}.$$

where $x_D = \Delta M_D/\Gamma_D$ and $y_D = \Delta \Gamma_D/2\Gamma_D$. Contours in the $(x_D,y_D)$ plane are shown in Fig. 2. The significance of the oscillation effect in the preliminary world averages exceeds $5\sigma$.

This result suggests that charm mixing may be at the upper end of the range of Standard Model predictions. The interpretation of this result in terms of New Physics is limited by the theoretical uncertainty on the Standard Model prediction. Nevertheless interesting constraints on the mixing amplitudes and NP contributions have been already deduced.

3 Phenomenological impact

The recent results have further improved the determination of the UT parameters. Assuming the SM, in Fig. 3 we show the results of the new fit which includes all constraints.

The numerical results are: $\bar{\rho} = 0.164 \pm 0.029$ ; $\bar{\eta} = 0.330 \pm 0.017$.

It is interesting to notice that the determination of the UT parameters from the measurements of CP violating quantities in the kaon $(\epsilon_K)$ and in the B sectors (UT angles) is now as precise as those obtained with the measurements of the sides (see Fig. 3).
Figure 2: Left plot: Babar result on $D$ mixing. a) Projections of the proper-time distribution of combined $D^+$ and $\bar{D}^0$ WS candidates and fit result integrated over the signal region. 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. Right plot: likelihood contours in the $(\tau, q_{13})$ plane as obtained by HFAG$^1$.

Figure 3: Determination of $\beta$ and $\bar{\eta}$. Left plot: constraints on $|V_{ud}|/|V_{us}|$, $\Delta m_d$, $\Delta m_s$, $|\Delta A|$, $\beta$, $\gamma$, and $\alpha$. 68% and 95% total probability contours are shown, together with 95% probability regions from the individual constraints. Right plot: The contours at 68%, 95% selected by the measurements of $|V_{ud}|/|V_{us}|$, $\Delta m_d$, and $\Delta m_s$ are compared to the bounds (at 95% probability) from the measurements of CP violating quantities in the $\tau$ and in the $B$ sectors.

In addition the recent measurements performed at $B$ factories and Tevatron allow for a simultaneous determination of the CKM parameters together with the NP contributions to $|\Delta F|=2$ processes. In fact each of the mixing processes is described by a single amplitude and can be parameterized, without loss of generality, in terms of two parameters, which quantify the difference of the complex amplitude with respect to the SM one$^{20}$. Thus, for instance, in the case of $B^0_q - \bar{B}^0_q$ mixing we define

$$C_{q_B} e^{2\delta_{q_B}} = \frac{|B^0_q|^{H_{	ext{full}}} |B^0_q|}{|B^0_q|^{H_{	ext{full}}} |B^0_q|}, \quad (q = d, s)$$

[1]
where \( H^{\text{SM}} \) includes only the SM box diagrams, while \( H^{\text{full}} \) includes also the NP contributions. In the absence of NP effects, \( C_{B_s} = 1 \) and \( \phi_{B_s} = 0 \) by definition.

The most important results are shown in Figs. 1-22. Two important conclusions can be drawn:

i) in the \( B_s \) sector we have already quite strong constraints on NP contributions which can be as large as the SM ones only if the SM and NP amplitudes have the same weak phase. More generally we can conclude that NP should contribute no more than 20\% with respect to the SM for a generic NP phase. In addition, the fit produces a nonzero central value for \( \phi_{B_s} \) due to the difference in the SM fit between the angles and the side measurements.

ii) The measurement of \( \Delta m_s \), strongly constrains \( C_{B_s} \) which is now better known than \( C_{B_d} \). \( \phi_{B_s} \) starts to be constrained. This result is coming from several new measurements: the semileptonic asymmetry in \( B_s \) decays, \( A_L \), the dimuon charge asymmetry (\( AC_H \)), the measurement of the \( B_s \) lifetime from flavon-specific final states, the determination of \( \Delta \Gamma_s / \Gamma_s \), from the time integrated analysis \( B_s \to J/\psi \phi \) and the three-dimensional constraint on \( \gamma_s \), the determination of \( \Delta \Gamma_s \) and the phase \( \phi_s \) of the \( B_s \) mixing amplitude from the time-dependent angular analysis of \( B_s \to J/\psi \phi \) [21,22].

4 Conclusions and perspectives

Flavour physics is a very active research field. Many interesting results are presented at this conference: new and more precise determination of the UT angles, of \( |V_{ub}| \), new results on radiative and purely leptonic decays which are very powerful instruments to look to NP... The highlight of the conference is the evidence of charm mixing. This result indicates that charm mixing is at the upper end of the range of Standard Model predictions and is already providing interesting constraints on the mixing amplitudes and NP contributions.

The Standard Model is still resisting against all these efforts to probe its validity limits. Nevertheless it has to be stressed that all the tests are at best at about 20\% level. The next facilities can surely push these tests to a 1\% or lower accuracy. In conclusions in Fig. 5 we show the regions on the \( \bar{p}-\bar{\eta} \) plane selected by different constraints assuming the current measurement and precision expected at high luminosity B-factory (Superb is a B-factory collecting 75ab\(^{-1}\)) [21]. With the precision reached at SuperB, the current discrepancies would clearly indicate the presence of New Physics in the flavour sector!
Figure 5: Regions corresponding to 95% probability for $\bar{\nu}$ and $\bar{\nu}$ selected by different constraints, assuming present central values with errors expected at SuperB.

Acknowledgements

I would like to warmly thanks the organisers of Moriond QCD for the kind invitation. Many thanks to Anne Marie Lutz for the careful reading of the document.

   Updates and plots: http://www.slac.stanford.edu/xorg/hfag/
5. B. Aubert et al. [BABAR Collaboration], [arXiv:hep-ex/0612021].


M. Ciuchini “Precision Flavour Physics” talk given at “Les Rencontres de
The study of charmless hadronic $B_s$ decays

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The perturbative QCD approach has achieved great success in the study of hadronic $B$ decays. Utilizing the constrained parameters in these well measured decay channels, we study most of the possible charmless $B_s \rightarrow PP, PV$ and $VV$ decay channels in the perturbative QCD approach. In addition to the branching ratios and CP asymmetries, we also give predictions to the polarization fractions of the vector meson final states. The size of SU(3) breaking effect is also discussed. All of these predictions can be tested by the future LHCb experiment.

Keywords: $B_s$ decays, factorization, pQCD

1 Introduction

There is a continuous progress in the study of hadronic $B$ decays since the so called naive factorization approach.\textsuperscript{1,2} In recent years, the QCD factorization approach (QCDF)\textsuperscript{3} and perturbative QCD factorization (pQCD) approach\textsuperscript{4} together with the soft-collinear effective theory\textsuperscript{5} solved a lot of problems in the non-leptonic decays. Although most of the branching ratios measured by the $B$ factory experiments can be explained by any of the theories, the direct CP asymmetries measured by the experiments are ever predicted with the right sign only by the pQCD approach.\textsuperscript{6} The LHCb experiment will soon run in the end of 2007. With a very large luminosity, it will accumulate a lot of $B_s$ events. The progress in both theory and experiment encourages us to apply the pQCD approach to the charmless $B_s$ decays in this work.\textsuperscript{7}

In the hadronic $B(B_s)$ decays, there are various energy scales involved. The factorization theorem allows us to calculate them separately. First, the physics from the electroweak scale down to b quark mass scale is described by the renormalization group running of the Wilson coefficients of effective four quark operators. Secondly, the hard scale from b quark mass scale to the factorization scale $\sqrt{\Lambda_{\text{QCD}}}$ are calculated by the hard part calculation in the perturbative QCD approach.\textsuperscript{5} When doing the integration of the momentum fraction $x$ of the light quark, end point singularity will appear in the collinear factorization (QCDF and SCET) which breaks down the factorization theorem. In the pQCD approach, we do not neglect the transverse momentum $k_T$ of the light quarks in meson. Therefore the endpoint singularity disappears. The inclusion of transverse momentum will also give large double logarithms $\ln^2 k_T$ and $\ln^2 x$ in the hard part calculations. Using the renormalization group equation, we can resum them for all loops to the leading order resulting Sudakov factors. The Sudakov factors suppress the endpoint contributions to make the calculation consistent.\textsuperscript{4}

The physics below the factorization scale is non-perturbative in nature, which is described by the hadronic wave functions of mesons. They are not perturbatively calculable, but universal for all the decay processes. Since many of the hadronic and semi-leptonic $B$ decays have been measured well in the two $B$ factory experiments, the light wave functions are strictly constrained. Therefore, it is useful to use the same wave functions in our $B_s$ decays determined from the hadronic $B$ decays. The uncertainty of the hadronic wave functions will come mainly from the SU(3) breaking effect between the $B_s$ wave function and $B$ wave function.\textsuperscript{7} In practice, we use a little larger $\omega_b$ parameter for the $B_s$ meson than the $B_d$.
Table 1: The branching ratios and CP asymmetry calculated in pQCD approach, QCDF and SCET approaches together with Experimental Data.

<table>
<thead>
<tr>
<th></th>
<th>SCET</th>
<th>QCDF</th>
<th>PQCD</th>
<th>EXP</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B(B_s \rightarrow K^\mp \pi^\pm)(10^{-6})$</td>
<td>4.9 $\pm$ 1.8</td>
<td>10 $\pm$ 6</td>
<td>11 $\pm$ 6</td>
<td>5.0 $\pm$ 1.3</td>
</tr>
<tr>
<td>$B(B_s \rightarrow K^- K^+)(10^{-6})$</td>
<td>18 $\pm$ 7</td>
<td>23 $\pm$ 27</td>
<td>17 $\pm$ 9</td>
<td>24 $\pm$ 5</td>
</tr>
<tr>
<td>$B(B_s \rightarrow \phi \phi)(10^{-6})$</td>
<td>22 $\pm$ 30</td>
<td>33 $\pm$ 13</td>
<td>14 $\pm$ 8</td>
<td></td>
</tr>
<tr>
<td>$A_{CP}(B_s \rightarrow K^\mp \pi^\pm)$ (%)</td>
<td>20 $\pm$ 26</td>
<td>$-6.7 \pm 16$</td>
<td>30 $\pm$ 6</td>
<td>39 $\pm$ 15 $\pm$ 8</td>
</tr>
</tbody>
</table>

meson, which characterize the fact that the light $s$ quark in $B_s$ meson carries a litter larger momentum fraction that the $d$ quark in the $B_s$ meson.

2 Results and Discussion

For $B_s$ meson decays with two light mesons in the final states, the light mesons obtain large momentum of 2.6GeV in the $B_s$ meson rest frame. All the quarks inside the light mesons are therefore energetic and collinear like. Since the heavy $b$ quark in $B_s$ meson carry most of the energy of $B_s$ meson, the light $s$ quark in $B_s$ meson is soft. In the usual emission diagram of $B_s$ decays, this quark goes to the final state meson without electroweak interaction with other quarks, which is called a spectator quark. Therefore there must be a connecting hard gluon to make it from soft like to collinear like. The hard part of the interaction becomes six quark operator rather than four. The soft dynamics here is included in the meson wave functions. The decay amplitude is infrared safe and can be factorized as the following formalism:

$$C(t) \times H(t) \times \Phi(x) \times \exp \left[ -s(P, b) - 2 \int_{\mu}^t \frac{d\mu}{\mu} \gamma_4(\alpha_s(\mu)) \right],$$

(1)

where $C(t)$ are the corresponding Wilson coefficients of four quark operators, $\Phi(x)$ are the meson wave functions and the variable $t$ denotes the largest energy scale of hard process $H$, which is the typical energy scale in pQCD approach and the Wilson coefficients are evolved to this scale. The exponential of $S$ function is the so-called Sudakov form factor resulting from the resummation of double logarithms occurred in the QCD loop corrections, which can suppress the contribution from the non-perturbative region. Since logarithm corrections have been summed by renormalization group equations, the above factorization formula does not depend on the renormalization scale $\mu$ explicitly.

The numerical results of the $B_s$ decays branching ratios and CP asymmetry parameters are displayed in ref.7. In all the decay channels for charmless $B_s$ decays, only several are measured by the CDF collaboration.9 We show those channels together with results by QCDF 10 and SCET approaches 11 in table 1. From those comparison, we notice that the measured branching ratios are still consistent with the theoretical calculations. Like the case in $B$ decays, the calculated branching ratios from the three kinds of methods overlap with each other, considering the still large theoretical uncertainties. A global fit is useful when we have enough measured channels.

In table 1, the only measured CP asymmetry in $B_s \rightarrow K^- \pi^+$ decay prefer our pQCD approach rather than QCDF approach. This is similar with the situation in $B$ decays. The direct CP asymmetry is proportional to the sine of the strong phase difference of two decay topologies.6 The strong phase in our pQCD approach is mainly from the chirally enhanced space-like penguin diagram, while in the QCDF approach, the strong phase mainly comes from the virtual charm quark loop diagrams. The different origin of strong phases gives different sign to the direct CP asymmetry imply a fact that the dominant strong phase in the charmless decays should come from the annihilation diagrams. It should be noted that the SCET approach can not predict the direct CP asymmetry of $B$ decays directly, since they need more experimental measurements as input. However, it also gives the right CP asymmetry for $B_s$ decay if with the input of experimental CP asymmetries of $B$ decays, which means good SU(3) symmetry here.

For the $B_s \rightarrow VV$ decays, we also give the polarization fractions in addition to the branching ratios and CP asymmetry parameters.12 Similar to the $B \rightarrow VV$ decay channels, we also have large transverse polarization fractions for the penguin dominant processes, such as $B_s \rightarrow \phi \phi$, $B_s \rightarrow K^{*+}K^{-}$, $K^{*0}K^{*0}$ decays, whose transverse polarization fraction can reach 40-50%.
3 SU(3) breaking effect

The SU(3) breaking effects come mainly from the $B_s(B_d)$ meson decay constant and distribution amplitude parameter, light meson decay constant and wave function difference, and various decay topology differences. As an example we mainly focus on the decays $B \to \pi \pi$, $B \to K \pi$, $B \to K \pi$ and $B \to K K$, as they can be related by SU(3)-symmetry. A question of considerable interest is the amount of SU(3)-breaking in various topologies (diagrams) contributing to these decays. For this purpose, we present in Table 2 the magnitude of the decay amplitudes (squared, in units of GeV$^2$) involving the distinct topologies for the four decays modes. The first two decays in this table are related by U-spin symmetry ($d \to s$) (likewise the two decays in the lower half). We note that the assumption of U-spin symmetry for the (dominant) tree ($T$) and penguin ($P$) amplitudes in the emission diagrams is quite good, it is less so in the other topologies, including the contributions from the $W$-exchange diagrams, denoted by $E$ for which there are non-zero contributions for the flavor-diagonal states $\pi^+\pi^-$ and $K^+K^-$ only. The U-spin breakings are large in the electroweak penguin induced amplitudes $P_{EW}$, and in the penguin annihilation amplitudes $P_{A}$ relating the decays $B_d \to K^+\pi^-$ and $B_s \to K^+K^-$. In the SM, however, the amplitudes $P_{EW}$ are negligibly small.

### Table 2: Contributions from the various topologies to the decay amplitudes (squared) for the four indicated decays. Here, $T$ is the contribution from the color favored emission diagrams; $P$ is the penguin contribution from the emission diagrams; $E$ is the contribution from the $W$-exchange diagrams; $P_{A}$ is the contribution from the penguin annihilation amplitudes; and $P_{EW}$ is the contribution from the electro-weak penguin induced amplitudes.

| mode $(\text{GeV}^2)$ | $|T|^2$ | $|P|^2$ | $|E|^2$ | $|P_{A}|^2$ | $|P_{EW}|^2$ |
|----------------------|--------|--------|--------|-------------|-------------|
| $B_d \to \pi^+\pi^-$ | 1.5    | 9.2 $\times$ 10$^{-3}$ | 6.4 $\times$ 10$^{-3}$ | 7.5 $\times$ 10$^{-3}$ | 2.7 $\times$ 10$^{-6}$ |
| $B_s \to \pi^+K^-$  | 1.4    | 7.4 $\times$ 10$^{-3}$ | 0      | 7.0 $\times$ 10$^{-3}$ | 5.4 $\times$ 10$^{-6}$ |
| $B_d \to K^+\pi^-$  | 2.2    | 18.8 $\times$ 10$^{-3}$ | 0      | 4.7 $\times$ 10$^{-3}$ | 7.4 $\times$ 10$^{-6}$ |
| $B_s \to K^+K^-$    | 2.0    | 14.7 $\times$ 10$^{-3}$ | 4.6 $\times$ 10$^{-3}$ | 9.8 $\times$ 10$^{-3}$ | 3.1 $\times$ 10$^{-6}$ |

In $B^0 \to K^-\pi^+$ and $\bar{B}^0 \to K^+\pi^-$, the branching ratios are very different from each other due to the differing strong and weak phases entering in the tree and penguin amplitudes. However, as shown by Gronau, the two relevant products of the CKM matrix elements entering in the expressions for the direct CP asymmetries in these decays are equal, and, as stressed by Lipkin subsequently, the final states in these decays are charge conjugates, and the strong interactions being charge conjugation invariant, the direct CP asymmetry in $B^0 \to K^-\pi^+$ can be related to the well-measured CP asymmetry in the decay $\bar{B}^0 \to K^+\pi^-$ using U-spin symmetry.

Following the suggestions in the literature, we can define the following two parameters:

\[
R_3 = \frac{|A(B_s \to \pi^+K^-)|^2 - |A(B_s \to \pi^-K^+)|^2}{|A(B_d \to \pi^-K^+)|^2 - |A(B_d \to \pi^+K^-)|^2},
\]

\[
\Delta = \frac{A_{CP}^2(B_d \to \pi^+K^-) - A_{CP}^2(B_s \to \pi^-K^+)}{A_{CP}^2(B_s \to \pi^-K^+) + BR(B_s \to \pi^+K^-) \tau(B_d)}.
\]

The standard model predicts $R_3 = -1$ and $\Delta = 0$ if we assume U-spin symmetry. Since we have a detailed dynamical theory to study the SU(3) (and U-spin) symmetry violation, we can check in pQCD approach how good quantitatively this symmetry is in the ratios $R_3$ and $\Delta$. We get $R_3 = -0.96_{-0.09}^{+0.11}$ and $\Delta = -0.03 \pm 0.08$. Thus, we find that these quantities are quite reliably calculable, as anticipated on theoretical grounds. SU(3) breaking and theoretical uncertainties are very small here, because most of the breaking effects and uncertainties are canceled due to the definition of $R_3$ and $\Delta$. On the experimental side, the results for $R_3$ and $\Delta$ are: \(^9\)

\[
R_3 = -0.84 \pm 0.42 \pm 0.15, \quad \Delta = 0.04 \pm 0.11 \pm 0.08.
\]

We conclude that SM is in good agreement with the data, as can also be seen in Fig. 1 where we plot theoretical predictions for $R_3$ vs. $\Delta$ and compare them with the current measurements of the same. The measurements of these quantities are rather imprecise at present, a situation which we hope will greatly improve at the LHCb experiment.
4 Summary

Based on the $k_T$ factorization, pQCD approach is infrared safe. Its predictions on the branching ratios and CP asymmetries of the $B^0(B^{+})$ decays are tested well by the B factory experiments. Using those tested parameters from these decays, we calculate a number of charmless decay channels $B_s \to PP$, $PV$ and $VV$ in the perturbative QCD approach. The experimental measurements of the three $B_s$ decay channels are consistent with our numerical results. Especially the measured direct CP asymmetry of $B_s \to \pi^-K^+$ agree with our calculations. We also discuss the SU(3) breaking effect in these decays, which is at least around 20-30%. We also show that the Gronau-Lipkin sum rule works quite well in the standard model, where the SU(3) breaking effects mainly cancel.

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References

Hadronic B decays at BELLE and BABAR

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There are several exciting results for the hadronic B decays of Belle and BaBar recently. My report focuses on the ratios of branching fractions and CP asymmetry for $B \to hh$ decays at Belle and BaBar where $h$ denotes $\pi$ or $K$. The observations of $B^+ \to K^+K^0$, $B^0 \to K^0\bar{K}^0$ are published both by Belle and BaBar and shown here. We also report the observation of $B^+ \to \rho^+ K^0$ and search of $B^+ \to \bar{K}^0 K^+$, $B^0 \to \rho^0 \rho^0$, $B^0 \to a_0 \pi^+$ and $B^0 \to a_0^- K^+$ at BaBar. Finally, we will show the results of amplitude analysis of the decays $B^0 \to \phi K^*_2(1430)^0$, $\phi K^*(892)^0$ and $\phi(K\pi)^0_{S-wave}$ at BaBar.

1 Introduction

In general, the branching fractions with the SM predictions suffer from large hadronic uncertainties within the current theoretical framework and many of the uncertainties cancel out in ratios of branching fractions. We report the ratios $R_{\epsilon_B}$ of the charged and neutral $B \to K\pi$ branching fractions and direct CP asymmetry of $B \to K\pi$. Theoretical predictions with different approaches suggest that $A_{CP}(K^+\pi^-)$ could be either positive or negative\(^1\). Although there are large uncertainties related to hadronic effects in the theoretical predictions, results for $A_{CP}(K^+\pi^-)$ and $A_{CP}(K^+\pi^0)$ are expected to have the same sign and be comparable in magnitude\(^1\). The observation of $B^0 \to K^0\bar{K}^0$ and $B^+ \to \bar{K}^0 K^+$ expected to be dominated by the loop-induced $b \to d \bar{s}s$ process (called a $b \to d$ penguins) are published by Belle and BaBar recently. BaBar also observed $B^+ \to \rho^+ K^0$ decay which is expected to be a pure penguin decays and the result helps to separate the contribution of tree and penguin amplitudes in other channels. For the amplitude analysis, we will show the results of $B^0 \to \rho^0 \rho^0$ and $B^0 \to \phi K^*_2(1430)^0$, $\phi K^*(892)^0$ and $\phi(K\pi)^0_{S-wave}$. The nature of the $a_0$ meson is still not well understand and the branching fractions of $B^0 \to a_0^- \pi^+$, $B^0 \to a_0^- K^+$, $B^0 \to a_0(1450)^-\pi^+$ and $B^0 \to a_0(1450)^-K^+$ will provide some information about the nature of $a_0$. 
Figure 1: $M_{bc}$ (left) and $\Delta E$ (right) distributions for $B^0 \rightarrow K^+\pi^-$, $B^+ \rightarrow K^+\pi^0$ and $B^+ \rightarrow \pi^+\pi^0_\text{c}$ candidates at Belle. The histograms show the data, while the curves represent the various components from the fit: signal (dotted), continuum (dashed), charmless $B$ decays (hatched), background from mis-identification (dotted), and sum of all components (solid). The $M_{bc}$ and $\Delta E$ projections of the fits are for events that have $|\Delta E| < 0.06$ GeV (left) and 5.271 GeV/c$^2 < M_{bc} < 5.289$ GeV/c$^2$ (right). (A looser requirement, $-0.14$ GeV $< \Delta E < 0.06$ GeV, is used for the modes with a $\pi^0_\text{c}$ meson in the final state.)

Figure 2: The background-subtracted distribution of $\Delta E$ for signal $K^+\pi^+$ events with data sample 347 M $B\bar{B}(left)$ and 383 M $B\bar{B}$ (right) at BaBar, comparing $B^0$ (solid) and $\bar{B}^0$ decays (dashed).

2 $B \rightarrow K\pi$, $\pi\pi$ and $KK$

The new measurements of the branching fractions for $B \rightarrow K^+\pi^-, K^+\pi^0, K^0\pi^0, \pi^+\pi^-,$ and $\pi^+\pi^0$ at Belle$^2$ and BaBar$^3,4$ are shown in Table 1. The effect of final-state radiation is considered in branching fraction measurement now. The statistical errors on the branching fraction for all decay modes are reduced. The ratio $R_c$ ($R_h$) obtained by Belle and BaBar’s experimental results are 1.08±0.06±0.08 (1.08±0.08±0.06) and 1.11±0.07 (0.94±0.07), respectively. The current $R_c$ and $R_h$ have moved quite a bit towards the SM predictions and reduce the “$B \rightarrow K\pi$ puzzle”.

The observation of $B^0 \rightarrow \bar{K}^0 K^0$ and $B^+ \rightarrow \bar{K}^0 K^+$ by Belle$^6$ and BaBar$^6$ are shown in Table 1 and the results agree with some theoretical predictions.$^7,8,9,10,11$

Throughout this letter, the partial-rate asymmetry is define as $A_{CP}(B \rightarrow f) = (\Gamma(B \rightarrow \bar{f}) - \Gamma(B \rightarrow f))/(\Gamma(B \rightarrow \bar{f}) + \Gamma(B \rightarrow f))$, where $\bar{B}$ and $\bar{f}$ are the conjugate states. Belle provides the partial-rate asymmetries for $B^0 \rightarrow K^+\pi^-$, $B^+ \rightarrow K^+\pi^0$ and $B^+ \rightarrow \pi^+\pi^0$ three decay modes and the results are shown in Fig. 1 and Table 2. The partial-rate asymmetry $A_{CP}(K^+\pi^-)$ is found to be $-0.093 \pm 0.018 \pm 0.008$, which $4.8\sigma$ from zero. The measurement of $A_{CP}(K^+\pi^0)$ is consistent with no asymmetry; the central value is $4.4\sigma$ away from $A_{CP}(K^+\pi^-)$. BaBar also provides the partial-rate asymmetry for these three decay modes shown in Fig. 2 and Table 2 and claims the observation of the partial-rate asymmetry $-0.107 \pm 0.018\pm0.008$ with $5.5\sigma$.$^{12}$
### Table 1: Summary of branching fractions

<table>
<thead>
<tr>
<th>Mode</th>
<th>$N_{\text{signal}}$</th>
<th>BF (10^{-6})</th>
<th>$N_{\text{signal}}$</th>
<th>BF (10^{-6})</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^+\pi^-$</td>
<td>3585\pm 198</td>
<td>19.9 \pm 0.4 \pm 0.8</td>
<td>1660\pm 52</td>
<td>19.7 \pm 0.6 \pm 0.6</td>
</tr>
<tr>
<td>$\pi^+\pi^-$</td>
<td>872 \pm 41</td>
<td>5.1 \pm 0.2 \pm 0.2</td>
<td>489 \pm 35</td>
<td>5.8 \pm 0.4 \pm 0.3</td>
</tr>
<tr>
<td>$K^+K^-$</td>
<td>2.5 \pm 0.5</td>
<td>&lt;0.41</td>
<td>3 \pm 13</td>
<td>&lt;0.4</td>
</tr>
<tr>
<td>$K^+\pi^0$</td>
<td>1493\pm 56</td>
<td>12.4 \pm 0.5 \pm 0.6</td>
<td>1239 \pm 52</td>
<td>13.3 \pm 0.6 \pm 0.6</td>
</tr>
<tr>
<td>$\pi^+\pi^0$</td>
<td>693\pm 39</td>
<td>6.5 \pm 0.4 \pm 0.4</td>
<td>572 \pm 53</td>
<td>5.1 \pm 0.5 \pm 0.3</td>
</tr>
<tr>
<td>$K^0\bar{K}^\pm$</td>
<td>36.6\pm 1.7</td>
<td>1.22 \pm 0.32 \pm 0.13</td>
<td>71 \pm 19</td>
<td>1.61 \pm 0.44 \pm 0.09</td>
</tr>
<tr>
<td>$K^0\pi^+$</td>
<td>1252\pm 41</td>
<td>22.8 \pm 1.0 \pm 0.6</td>
<td>1072 \pm 46</td>
<td>23.9 \pm 1.1 \pm 1.0</td>
</tr>
<tr>
<td>$K^{0}\bar{K}^0$</td>
<td>23.0\pm 1.6</td>
<td>0.87 \pm 0.20 \pm 0.09</td>
<td>32 \pm 8</td>
<td>1.08 \pm 0.28 \pm 0.11</td>
</tr>
<tr>
<td>$K^{0}\pi^0$</td>
<td>379\pm 8</td>
<td>9.2 \pm 0.7 \pm 0.6</td>
<td>425 \pm 28</td>
<td>10.5 \pm 0.7 \pm 0.5</td>
</tr>
<tr>
<td>$\rho^+\bar{K}^0$</td>
<td>-</td>
<td>-</td>
<td>158\pm 27</td>
<td>8.0\pm 1.4 \pm 0.5</td>
</tr>
</tbody>
</table>

### Table 2: Summary of partial-rate asymmetry

<table>
<thead>
<tr>
<th>Mode</th>
<th>Belle(535M $BB$)</th>
<th>BaBar(347M $BB$)</th>
<th>BaBar(383M $BB$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$K^+\pi^-$</td>
<td>-0.093 \pm 0.018 \pm 0.008</td>
<td>-0.108 \pm 0.024 \pm 0.007</td>
<td>-0.107 \pm 0.018 \pm 0.007</td>
</tr>
<tr>
<td>$K^+\pi^0$</td>
<td>0.07 \pm 0.03 \pm 0.01</td>
<td>0.016 \pm 0.041 \pm 0.010</td>
<td>-</td>
</tr>
<tr>
<td>$\pi^+\pi^0$</td>
<td>0.07 \pm 0.06 \pm 0.01</td>
<td>-0.019 \pm 0.088 \pm 0.014</td>
<td>-</td>
</tr>
</tbody>
</table>

3. $B^+ \to \rho^+K^0$, $B^0 \to \rho^0\bar{\rho}^0$ and $B^0 \to \phi K^{*0}$

The pure penguin $b \to s$ decay process $B^+ \to \rho^+K^0$ is observed by BaBar recently\,\,14. The measured branching fraction is $(8.0^{+1.4}_{-1.3} \pm 0.5) \times 10^{-6}$ shown in Table 1 with 7.9\,$\sigma$ significance and is consistent with theoretical prediction with the assumption $p_{\rho^+} = -p_{\rho^-}$\,\,13 within the uncertainties. The $p_{\rho^+}$ ($p_{\rho^-}$) is the amplitude for the spectator quark to appear in the vector (pseudoscalar) meson.

The value of $\Delta\alpha$ can be extracted from an analysis of the branching fractions of the $B$ decays into the full set of isospin-related channels\,\,15. BaBar finds the evidence for $B^0 \to \rho^0\bar{\rho}^0$ with 3.5\,$\sigma$ significance and measures the branching fraction $(1.07 \pm 0.33 \pm 0.19) \times 10^{-6}$\,\,16. With the $B^0 \to \rho^0\bar{\rho}^0$ measurement, BaBar obtains a 68\% (90\%) CL limit on $|\Delta\alpha| \equiv |\alpha - \alpha_{\text{eff}}| < 18^\circ$ (< 20\(^\circ\)). An isospin-triangle relation holds for each of the three helicity amplitudes, which can be separated through an angular analysis. The longitudinal polarization fraction $f_L = |A_0|^2/(\Sigma |A_\lambda|^2)$ of $\rho^0\bar{\rho}^0$ is $0.87 \pm 0.13 \pm 0.04$, where $A_{\lambda = -1,0,+1}$ are the helicity amplitudes.

The large fraction of transverse polarization in the $B \to \phi K^{*}(892)$ decay measured by Belle\,\,17 and BaBar\,\,18 indicates a significant departure from the naive expectation of dominant longitudinal polarization. BaBar extend their investigation of the polarization puzzle with an amplitude analysis of the vector-tensor $B^0 \to \phi K_T^0(1430)^0$ decay and vector-scalar $B^0 \to \phi(K\pi)^0$ decay\,\,19. The amplitudes are reparameterized with the index $J$ suppressed as $A_D$ and $A_{\pm 1} = (A_{\pm 1})/\sqrt{2}$ and the transverse polarization fraction is defined as $f_T = |A_{\pm 1}|^2/\Sigma |A_\lambda|^2$. The polarization results are

$$f_L(B^0 \to \phi K^{*}(892)^0) = 0.506 \pm 0.040 \pm 0.015$$
$$f_L(B^0 \to \phi K_T^0(1430)^0) = 0.853^{+0.061}_{-0.065} \pm 0.036$$
$$f_L(B^0 \to \phi K^{*}(892)^0) = 0.227 \pm 0.038 \pm 0.013$$
\[ f_\perp (B^0 \to \phi K^*_2(1430)^0) = 0.045^{+0.049}_{-0.040} \pm 0.013 \]

4. \( B \to a_0 K \) and \( a_0 \pi \)

BaBar apply separate fits to determine the \( a_0(980) \) and \( a_0(1450) \) yields since this results in \( \sim 20% \) better sensitivity for \( a_0(980) \). The \( a_0(1450) \) fit has a component for \( a_0(980) \) with the yiled fixed to the value found in the \( a_0(980) \) fit, corrected for the small efficiency difference. Since the branching fraction for \( a_0 \to \eta \pi \) is not well known, the following is the 90\% C.L. upper limits of product branching fraction\(^{20}\). Since the branching fractions of these decays are similar, it means the \( a_0(980) \) meson tend to four-quark state\(^{21}\).

\[
\begin{align*}
\mathcal{B}(B^0 \to a_0^- \pi^+) \times (a_0 \to \eta \pi) & < 3.1 \times 10^{-6} \\
\mathcal{B}(B^0 \to a_0^- K^+) \times (a_0 \to \eta \pi) & < 1.9 \times 10^{-6} \\
\mathcal{B}(B^0 \to a_0(1450)^- \pi^+) \times (a_0 \to \eta \pi) & < 2.3 \times 10^{-6} \\
\mathcal{B}(B^0 \to a_0(1450)^- K^+) \times (a_0 \to \eta \pi) & < 3.1 \times 10^{-6}
\end{align*}
\]

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16. B. Aubert et al. (BaBar Collaboration), hep-ex/0612021.
20. B. Aubert et al. (BaBar Collaboration), hep-ex/0703038.
Top Quark Mass Measurements at the Tevatron and the Standard Model Fits

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New measurements of the top quark mass from the Tevatron are presented. Combined with previous results, they yield a preliminary new world average of $m_{\text{top}} = 170.9 \pm 1.1 \text{(stat)} \pm 1.5 \text{(syst)} \text{GeV}/c^2$ and impose new constraints on the mass of the Higgs boson.

1 Introduction

The huge interest in a precise measurement of the top quark mass ($m_{\text{top}}$) is primarily motivated by its role in constraining the mass of the Higgs boson ($m_{\text{Higgs}}$). To see this, let us begin by looking at the mass of the $W$ boson ($m_W$) in the Standard Model which, when one-loop radiative corrections are included, can be related to well known electroweak quantities through the following expression:

$$m_W^2 = \frac{m_0^2}{\sin^2 \theta_W (1 + \Delta r)}.$$  \hspace{1cm} (1)

The radiative corrections contained in $\Delta r$ receive contributions from the top quark:

$$(\Delta r)_{\text{top}} \approx \frac{3G_F m_{\text{top}}^2}{8\sqrt{2}\pi^2} \frac{1}{\tan^2 \theta_W}$$  \hspace{1cm} (2)

and the Higgs boson:

$$(\Delta r)_{\text{Higgs}} \approx \frac{11G_F m_Z^2 \cos^2 \theta_W}{24\sqrt{2}\pi^2} \ln \frac{m_{\text{Higgs}}^2}{m_Z^2}$$  \hspace{1cm} (3)

where $m_Z$ is the mass of the $Z$ boson. From these expressions, we see that $m_{\text{top}}$ enters quadratically while $m_{\text{Higgs}}$ enters logarithmically. A precise knowledge of both $m_W$ and $m_{\text{top}}$ in combination with existing electroweak data is therefore necessary to impose useful constraints on $m_{\text{Higgs}}$. Such constraints, in turn, are of tremendous value in the ongoing search for the Higgs.

In this talk, we present the latest top quark mass measurements from the CDF and DØ collaborations based on up to 1 fb$^{-1}$ of Run II data collected at Fermilab's Tevatron. These results are combined with previous ones to give a new preliminary world average for $m_{\text{top}}$ which, in turn, yields new constraints on the Higgs mass.
2 Measurement Channels and Experimental Challenge

Now that we understand the motivation behind a precise determination of the top mass, let us look at the top quark decay channels in which these measurements are performed and the experimental challenges they pose.

In the all jets channel, both $W$ bosons from the $t\bar{t}$ pair decay hadronically into jets for a total of 6 jets in the event. This channel has the advantage of having the largest branching ratio of 44%. It suffers, however, from large background levels from QCD multijet events. On the other hand, it benefits from the presence of the hadronically decaying $W$ bosons whose well known masses can be exploited to perform an in-situ calibration of the jet energies, reducing the effect of the systematic uncertainty in the overall jet energy scale. In the dilepton channel, both $W$ bosons decay leptonically. It has the advantage of having the lowest background levels coming from Drell-Yan processes associated with jets, diboson production with associated jets, and $W$+3 jet events with one jet faking an electron. Unfortunately, it also has the lowest branching ratio of 5%. In the lepton+jets ($l$+jets) channel, one of the two $W$ bosons from the $t\bar{t}$ pair decays hadronically while the other one decays leptonically. This channel maintains a good balance between a reasonable branching ratio of 29% and moderate background levels from $W$+jets and QCD multijet events. Like the all jets channel, it can benefit from an in-situ jet energy calibration using the $m_T$ constraint. It has traditionally yielded the most precise $m_{top}$ measurements.

To appreciate the challenge involved in measuring the top mass at the Tevatron, let us now take the $l$+jets channel as an example. In this case, what our reconstruction programs give us from the detector are several jets, a high $p_T$ lepton, substantial missing transverse energy, and an interaction vertex. Since we don’t really know how to associate jets with partons in general, all jet permutations need to be considered in a straightforward reconstruction of the top mass. Furthermore, unlike long lived particles, there are no detached vertices associated with the top quark itself that can be used to separate the signal from the background events. This means that, even with $b$-tagging, there are no sharp and clean mass peaks from which the top mass can be determined directly. Fortunately, despite these challenges, sophisticated measurement techniques have been developed that make a precise measurement of the top mass possible.

3 Top Quark Mass Measurement Techniques

In this section we describe the three major techniques used in measuring the top quark mass. All the measurements presented here use one or some combination of these techniques.

The template method is the oldest of the three techniques and has been used for most of the earliest mass measurements. In this technique, one begins by identifying a variable sensitive to the top mass, an obvious choice of which would be the kinematically reconstructed value of the mass itself. Distributions of the chosen variable are then plotted separately for several samples of fully simulated Monte Carlo (MC) events differing only in the value of the top mass used to generate the signal events. Each of these distributions is called a template and is associated with a particular value of the input mass. The top mass is then extracted from the data sample by comparing the data distribution directly with each MC template to find the best fit value based on some measure of the goodness of fit. More recent applications of this technique parameterize the templates in terms of a probability density function which is used to construct likelihoods from which the top mass is extracted.

DØ pioneered the application of the matrix element (ME) method to top quark mass measurements in the Run I data from the $l$+jets channel. It is based on calculating the probability for observing each event which includes contributions from both signal and background sources. The signal probability is calculated as a function of the assumed top mass, resulting in a prob-
Figure 1: 2D likelihoods for electron and muon channels for the DØ $\ell$+jets result.

ability distribution for each event. The probability is taken to be the differential cross section for the process in question. The calculated probability distributions for every event in the data sample are combined to construct a joint likelihood from which the top mass is determined and its uncertainty estimated. The ME method makes use of as many measured variables as possible to completely specify an event, thereby allowing maximum discrimination between signal and background events. Within each event, all possible jet permutations are combined in a natural way based on their relative probabilities. Furthermore, the use of transfer functions allows a probabilistic treatment of the mapping between parton and jet energies where the full spectrum of parton energies contributing to the observed jet energy is taken into account.

The ideogram method, like the ME method, calculates an event-by-event likelihood. This technique makes use of a constrained kinematic fit to reconstruct the top mass. Using a simple parameterization, the probability for observing the reconstructed mass is then calculated as a function of the true value with the measurement resolution taken into account. This technique, which was also pioneered by DØ\textsuperscript{2}, aims to achieve statistical uncertainties comparable to those of the ME method without requiring as many computational resources.

4 New Results from the Tevatron

DØ has measured the top quark mass in the $\ell$+jets channel using the ME method described in the previous section\textsuperscript{3}. This measurement takes advantage of the $m_{V}$ constraint to perform an in-situ calibration of the jet energies. This is done by introducing a global scale factor, $JES$, that is applied to the energies of all the jets. A fit is then performed that maximizes the likelihood simultaneously in $m_{top}$, $JES$, and the signal fraction $C_{s}$. The 2D likelihood fits in $m_{top}$ and $JES$ are shown separately for the electron and muon channels in Figure 1. The combined result for both channels is $170.5 \pm 2.4$(stat + JES) $\pm 1.2$(syst)GeV/c$^2$ for 0.9 fb$^{-1}$ of data. Dominant systematic uncertainties are in the modeling of initial and final state radiations and $b$-fragmentation. This is the best DØ measurement of the top quark mass to date.

CDF has also measured the top quark mass in the $\ell$+jets channel using the ME method. Like the DØ result, this measurement employs an in-situ jet energy calibration through the inclusion of a global $JES$ parameter in the likelihood fit. The left plot in Figure 2 shows the 2D likelihood fit to the data for both electron and muon channels in $JES$ and $m_{top}$. The right plot in Figure 2 shows the expected error distribution from MC ensemble tests with the arrow indicating the measurement uncertainty. The measured result for 0.94 fb$^{-1}$ of data is $170.9\pm 2.2$(stat + JES) $\pm 1.4$(syst)GeV/c$^2$. The largest systematic uncertainty is in the modeling of initial and final state radiations. This is currently the most precise CDF measurement of the
top quark mass.

DØ has a measurement of the top quark mass in the $\ell+$jets channel using the ideogram method\textsuperscript{5}. Like the two results above, it employs an in-situ jet energy calibration. The 2D likelihood as a function of $JES$ and $m_{\text{top}}$ is shown on the left in Figure 3 with the gray line indicating the fitted value of $JES$ as a function of $m_{\text{top}}$. The right plot in Figure 3 shows the 1D likelihood as a function of $m_{\text{top}}$ along the gray line in the left plot. The result for 0.4 fb$^{-1}$ of data is $173.7 \pm 4.4(\text{stat} + \text{JES})^{+2.1}_{-1.0}(\text{syst})\text{GeV}/c^2$. Dominant systematic uncertainties are in the modeling of $b$-fragmentation and in the $b$/light jet energy scale ratio.

CDF has applied the ME method to a measurement of the quark top mass in the dilepton channel\textsuperscript{5}. A plot of the probability as a function of $m_{\text{top}}$ is shown on the left in Figure 4 and the expected error distribution on the right with the arrow indicating the measurement uncertainty. The result for 1 fb$^{-1}$ of data is $164.5 \pm 3.9(\text{stat}) \pm 3.9(\text{syst})\text{GeV}/c^2$. The systematic error is dominated by the uncertainty in the jet energy scale.

DØ has measured the top quark mass in the dilepton channel using a template method that assigns a weight to each neutrino solution based on the agreement between the calculated transverse momentum of the neutrinos and the observed missing transverse energy\textsuperscript{6}. The result for 1 fb$^{-1}$ is $172.5 \pm 5.8(\text{stat}) \pm 5.5(\text{syst})\text{GeV}/c^2$. The dominant source of the systematic error is the jet energy scale uncertainty.

CDF has measured the top quark mass in the all jets channel using a combination of template and ME methods\textsuperscript{7}. Instead of using the ME method directly to measure the top mass, the value
determined from the method is used to construct the MC templates. Probabilities calculated from the ME are also used in the event selection process to identify events with high signal probability. This result also uses the $m_W$ constraint to perform an in-situ jet energy calibration. The left plot in Figure 5 shows a fit of the data distribution to the MC templates for events with two b-tagged jets. Contours of JES and $m_{top}$ in data are shown on the right in Figure 5. The result for 1 fb$^{-1}$ is $171.1 \pm 3.7$ (stat + JES) $\pm 2.1$ (syst) GeV/c$^2$. The largest systematic uncertainties are in the simulation of fragmentation and showering and of final state radiation.

5 New World Average and Standard Model Fits

From above, the best result of each experiment in each channel is combined with previous results yielding a new preliminary world average$^8$ of $m_{top} = 170.9 \pm 1.1$ (stat) $\pm 1.5$ (syst) GeV/c$^2$ shown on the left in Figure 6. The ME $\ell$+jets results from DO and CDF carry the largest weights in this average of 40% and 39%, respectively. This value is 0.5 GeV/c$^2$ lower than the previous world average. With this new preliminary result, the top quark mass is now known to a total uncertainty of 1.8 GeV/c$^2$ corresponding to a relative precision of 1.1%.

This new top quark mass is also combined with other precision electroweak results in Standard Model fits performed by the LEP Electroweak Working Group$^9$. The right plot in Figure 6 shows the $\Delta \chi^2$ curve resulting from these fits giving $m_{Higgs} = 76^{+33}_{-24}$ GeV/c$^2$ at the minimum and a 95% confidence level upper limit of 144 GeV/c$^2$ which increases to 182 GeV/c$^2$ when the
LEP-2 direct search limit of 114 GeV/c² indicated by the yellow band is included.

6 Summary and Conclusions

A precise determination of the top quark mass is crucial for constraining the mass of the Higgs boson. Despite the great challenges involved, precise measurements are possible through the use of sophisticated measurement techniques. This talk presented new results based on up to 1 fb⁻¹ of data collected by CDF and DO. Although these results are still dominated by the $\ell+\text{jets}$ channel, the other two show promise and we hope to see more competitive results from them in the future. Combining the new results with previous ones has yielded a new preliminary world average top quark mass with a total uncertainty of 1.8 GeV/c² and imposed new constraints on $m_{\text{Higgs}}$. As more data become available at the Tevatron, we can expect statistical uncertainties < 1 GeV/c² by the end of the Tevatron run at which point the total uncertainties will become dominated by the systematic uncertainties.

References

MEASUREMENT OF MASSES AND LIFETIMES OF B HADRONS

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We present recent measurements by the CDF and DØ Collaborations at the Tevatron Collider on the masses and lifetimes of B hadrons. The results are compared to predictions based on Heavy Quark Effective Theory, lattice gauge theory, and quark models.

Keywords: hadrons, spectroscopy, lifetimes, Tevatron

1 Introduction

The ongoing Run II at the Tevatron pp Collider at the Fermi National Accelerator Laboratory has produced a wealth of results on B-physics. While the physics of $B_d$ and $B^\pm$ mesons is largely the domain of the $e^+e^-$ B-factories operating at the $\Upsilon(4S)$ resonance, access to the heavier B hadrons is presently reserved exclusively for the Tevatron. The drawback of working in the less clean hadron collider environment is more than compensated for by the high total $b\bar{b}$ production cross section (of order 100 $\mu$b) and the high luminosity (the integrated luminosity delivered to the CDF and DØ experiments to date is of order 2 $fb^{-1}$).

In the following, we discuss recent results obtained thanks to the high statistics B samples. Sect. 2 discusses new excited states that have been identified in the $B$, $B_s$, and $\Lambda_b$ systems. Sect. 3 covers a number of new and precise lifetime measurements. The measurements have been made on data samples of integrated luminosity ranging from 0.3 to 1.3 $fb^{-1}$. The two experiments’ results are compared, both with each other and theoretical predictions.

2 Spectroscopy

2.1 Orbitally excited $B$ and $B_s$ mesons

The spectroscopy of bound states containing one heavy quark ($Q$) and a light antiquark$^c$ is of great interest to quark models, as the dynamics of the light antiquark becomes independent of $m_Q$, and the total and light quark’s angular momentum become good independent quantum numbers. However, low statistics prevented the $L = 1$ orbitally excited $B^{**}$ states from being investigated in great detail previously.

The $B^{**}$ system is thought to consist of four states, two of which have a light quark angular momentum $J_q = 1/2$; their decay to $B^{(*)}\pi$ proceeds via an S-wave, and their total decay width is expected to be $\mathcal{O}(100\text{MeV})$, too wide to be observed unambiguously. The $J_q = 3/2$ $B_1$ and

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*aHere and in the following, charge conjugated states are implied.
B_{2}^{s}$ states (with total spin 1 and 2, respectively) decay via a $D$-wave, and their width, of order 10 MeV, should allow them to be identified clearly.

Both experiments have searched for the $B_{1}^{0}$ and $B_{2}^{0}$ states, through the decays $B_{1}^{0} \rightarrow B^{+}\pi^{-}$ and $B_{2}^{0} \rightarrow B^{(*)}\pi^{-}$. The photon from the decay $B^{(*)} \rightarrow B^{+}\gamma$ is not observed, resulting in two $B_{2}^{0}$ peaks in the $B\pi$ invariant mass spectrum, displaced by 46 MeV. Both experiments identified $B^{+}$ mesons through their decay to $J/\psi K^{+}$, the $J/\psi$ being recognized easily in its decay to muons; in addition, CDF employed the decay $B^{+} \rightarrow \bar{D}^{0}\pi^{+}$.

The resulting $B\pi$ invariant mass spectra, for the $J/\psi K^{+}$ decay mode, are shown in Fig. 1. CDF’s results are consistent with those using the $D^{*}\pi^{+}$ mode. The two experiments’ results obviously disagree: while DØ find clear evidence for three separately observed peaks, CDF find that the $B_{1}^{0}$, $B_{2}^{0} \rightarrow B^{(*)}\pi^{-}$ peaks coincide. The resulting $B_{1}^{0}$–$B_{2}^{0}$ mass splittings (25 MeV and 4 MeV, respectively) both disagree with the theory prediction of 12–14 MeV.

Both experiments have also carried out analogous analyses of $B_{s}^{(*)0} \rightarrow B^{+}\pi^{-}$ decays. The $B_{s}^{(*)}$ system is expected to replicate that of the $B^{(*)}$. However, the decay $B_{s}^{(*)0} \rightarrow B^{(*)}\pi^{-}$ is expected to be suppressed strongly, so that at most two peaks should be discernible. The results obtained are shown in Fig. 2. Both experiments find a very clear peak around 67 MeV, which is attributed to $B_{s}^{(*)0} \rightarrow B^{(*)}\pi^{-}$ decays. In addition, CDF find evidence for $B_{s}^{(*)0} \rightarrow B^{(*)}\pi^{-}$ decays, at a $Q$-value of 11 MeV (DØ would not expect to observe a $B_{s}^{(*)0}$ peak, given their observed mass splitting in the $B$ system). The resulting mass splitting of 10 MeV is fairly consistent with theoretical predictions.
Summarizing, both experiments obtain results consistent between B and B_s; however, there is a lack of agreement between CDF and DØ.

### 2.2 Σ_b baryons

CDF have used a sample of fully reconstructed $Λ_b^0 → Λ^+_c π^-$ decays to search for $Σ_b$ baryons, in the decay $Σ_b^{(*)±} → Λ^0_π^±$. In these baryons, the light di-quark system has $I = 1$ and $J^P = 1^+$; adding the b quark leads to $J^P = \frac{1}{2}^+$ ($Σ_b^0$) and $J^P = \frac{3}{2}^+$ ($Σ_b^{±}$). The $Σ_b – Λ_b$ mass difference is expected to be of order 200 MeV. The mass splitting between the two systems is expected to be $m(Σ_b^0) - m(Σ_b^+) = m_c/m_b \cdot (m(Σ_b^0) - m(Σ_b^±)) ≈ 20$ MeV; in addition, a small splitting between states within the same isospin triplet is expected.

In the analysis, the total decay widths are assumed to be saturated by single pion transitions, and constrained to the theory expectation of ~ 8 MeV and 15 MeV for the $Σ_b$ and $Σ_b^{±}$, respectively. In addition, the isospin splittings are assumed to be the same for $Σ_b$ and $Σ_b^+$. The result is shown in Fig. 3, and is in good agreement with theory predictions: $m(Σ_b^0) = 5808^{+2.3}_{-2.0} ± 1.7$ MeV, $m(Σ_b^+) = 5808^{+2.3}_{-2.0} ± 1.7$ MeV, $m(Σ_b^{±}) = 5837^{+2.1}_{-1.9} ± 1.7$ MeV, $m(Σ_b^{±}) = 5829^{+2.0}_{-1.8} ± 1.7$ MeV. (Note that not all of these results are independent.)

![Figure 3: Distribution of the mass difference $m(Λ_b^0π) - m(Λ_b^0) - m(π)$ for same-charge and opposite-charge $Σ_b$ signals.](image)

### 2.3 The $B_c$ meson

The $B_c$ meson, the lowest mass bound state of a b and c quark, has been observed already at LEP and in Run I of the Tevatron. However, only two fully reconstructed event candidates were observed so far, in the decay mode $B_c^+ → J/ψπ^+$. The resulting uncertainty on the mass was relatively large, about 60 MeV. CDF have now used this same decay mode in their Run II data sample, requiring in addition that the $J/ψ$ decay vertex be significantly displaced from the interaction point. They observe a signal of 49.1 ± 9.7 events over a background of 34.1 events. The result of the mass fit is $m(B_c) = 6276.5 ± 4.0$(stat.) ± 2.7(syst.) MeV. This represents an improvement of an order of magnitude in accuracy over previous results.

### 3 Lifetimes

Heavy Quark Effective Theory allows for a systematic expansion in orders of $α_s$ and $1/m_Q$ of the total decay widths of heavy-quark hadrons. As a result, precise predictions have been made for the ratios of lifetimes of B hadrons. In the past, this led to the so-called “$Λ_b$ puzzle”, where the measured $Λ_b$ lifetime ratio $τ(Λ_b)/τ(B_d)$ was significantly below its theoretical expectation. Improved lattice gauge theory computations have decreased this theoretical expectation substantially to $0.88 ± 0.05$, in fair agreement with experiment.
Both CDF and DØ have measured $\tau(A_b)$ in the exclusive decay $A_b \rightarrow J/\psi(\rightarrow \mu^+\mu^-)\Lambda(\rightarrow p\pi)$. This decay is very similar to the decay $B_d \rightarrow J/\psi K_S(\rightarrow \pi^+\pi^-)$, allowing for a “calibration” of the analyses using the precisely known $B_d$ lifetime. The $J/\psi$ decay vertex is combined with the reconstructed $\Lambda$ ($K_S$) track to yield the $A_b$ ($B_d$) vertex in the plane perpendicular to the beam line; correcting for the boost yields the lifetime estimate. Simultaneous unbinned maximum likelihood fits were made to the invariant mass and lifetime distributions (and given that event-by-event lifetime resolution estimates are used, the lifetime resolution distribution).

Both experiments measure a $B_d$ lifetime compatible with the world average. For the $A_b$, DØ’s measurement, $\tau(A_b)/\tau(B_d) = 0.811_{-0.038}^{+0.085}\text{(stat.)} \pm 0.034\text{(syst.)}$ is compatible with previous measurements and the theoretical predictions. However, the purer and more precise CDF result, $\tau(A_b)/\tau(B_d) = 1.018 \pm 0.062\text{(stat.)} \pm 0.007\text{(syst.)}$, is much higher than previous estimates.

DØ have measured the same quantity in the inclusive semileptonic decay $A_b \rightarrow \Lambda_c^+\mu^-\bar{\nu}_\mu X$. This measurement differs in many respects: it uses large statistics, but as it is an inclusive measurement a $A_b$ peak is not observed. Instead, the $\Lambda_c$ is reconstructed in its decay mode $\Lambda_c \rightarrow Ksp$. It is observed on top of a large background, which is subtracted statistically, in bins of the visible proper decay length. This quantity is corrected for the particles not reconstructed in the $\Lambda_c$ decay, the correction being modeled by Monte Carlo simulations. The result, $\tau(\Lambda_c) = 1.28_{-0.12}^{+0.13}\text{(stat.)} \pm 0.09\text{(syst.)} \text{ps}$, is in good agreement with DØ’s measurement in the $J/\psi\Lambda$ channel. In conclusion, the “$A_b$ puzzle” cannot yet be considered resolved – but at least we know there is an experimental issue to be addressed!

CDF have used their $J/\psi$ sample to measure also other lifetimes in exclusively reconstructed decays. In particular, using the decays $B^+ \rightarrow J/\psi K^+$ and $B^0 \rightarrow J/\psi\phi, \phi \rightarrow K^+K^-$, they find $\tau(B^+) = 1.630 \pm 0.016\text{(stat.)} \pm 0.011\text{(syst.)} \text{ps}$ and $\tau(B^0) = 1.494 \pm 0.054\text{(stat.)} \pm 0.009\text{(syst.)} \text{ps}$, respectively. These measurements are in good agreement with earlier measurements, as well as with HQET predictions. Similarly, a new DØ measurement of the $B_s$ lifetime using $B^0 \rightarrow D_s^-\mu^+\nu_\mu X$ decays, with the $D_s^-$ identified through $D_s^- \rightarrow \phi\pi^-, \phi \rightarrow K^+K^-$, yields $\tau(B^0_s) = 1.398 \pm 0.044\text{(stat.)}^{+0.038}_{-0.025}\text{(syst.)} \text{ps}$. This is in fair agreement with the current world average result. It should be pointed out that both experiments’ new $B_s$ results have accuracies comparable to the present world average, offering good hopes for more incisive tests of HQET.

Finally, CDF carried out a partial reconstruction of the decay $B_c^+ \rightarrow J/\psi(\rightarrow \mu^+\mu^-)e^+\nu_e X$. The $B_c$ lifetime is expected to be much shorter, $\tau(B_c) = 0.48 \pm 0.05 \text{ps}$, than that of the other weakly decaying $B$ hadrons, because also the charm quark can decay. The analysis attempts to isolate the exclusive decay $B_c^+ \rightarrow J/\psi e^+\nu_e$ using tight kinematic cuts. The challenge is a proper understanding the instrumental backgrounds (many of which are non-prompt), inferred from data, and from $b\bar{b}$ production, estimated using MC. The result, $\tau(B_c) = 0.463_{-0.036}^{+0.079}\text{(stat.)} \pm 0.036\text{(syst.)} \text{ps}$, is in good agreement with the present world average, $\tau(B_c) = 0.45_{-0.18}^{+0.19}\text{(stat.)} \pm 0.03\text{(syst.)} \text{ps}$. The new result is the most precise one to date, and its accuracy approaches that of the theoretical predictions.

References

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PRECISE CHARM AND BOTTOM QUARK MASSES

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New data for the total cross section $\sigma(e^+e^- \rightarrow \text{hadrons})$ in the charm and bottom threshold region are combined with an improved theoretical analysis, which includes recent four-loop calculations, to determine the short distance MS charm and bottom quark masses. The final result for the MS-masses, $m_c(3\,\text{GeV}) = 0.986(13)\,\text{GeV}$ and $m_b(10\,\text{GeV}) = 3.609(25)\,\text{GeV}$ is consistent with but significantly more precise than a similar previous study.

The strong coupling constant and the quark masses are the fundamental input parameters of the theory of strong interaction. Quark masses are an essential input for the evaluation of weak decay rates of heavy mesons and for quarkonium spectroscopy. Decay rates and branching ratios of a light Higgs boson, suggested by electroweak precision measurements, depend critically on the masses of the charm and bottom quarks, $m_c$ and $m_b$. Last not least, confronting the predictions for these masses with experiment is an important task for all variants of Grand Unified Theories. To deduce the values in a consistent way from different experimental investigations and with utmost precision is thus a must for current phenomenology.

A detailed analysis of $m_c$ and $m_b$ based on the ITEP sum rules\textsuperscript{1} has been performed several years ago\textsuperscript{2} and lead to $m_c(m_c) = 1.304(27)\,\text{GeV}$ and $m_b(m_b) = 4.191(51)\,\text{GeV}$. During the past years new and more precise data for $\sigma(e^+e^- \rightarrow \text{hadrons})$ have become available in the low energy region, in particular for the parameters of the charmonium and bottomonium resonances. Furthermore, the error in the strong coupling constant $\alpha_s(M_Z) = 0.1189 \pm 0.0020$ has been reduced. Last not least, the vacuum polarization induced by massive quarks has recently been computed in four-loop approximation\textsuperscript{3,4}; more precisely: its first derivative at $q^2 = 0$ has been evaluated, which corresponds to the lowest moment of the familiar $R$-ratio. With the help of the traditional integration-by-parts method in combination with Laporta’s algorithm\textsuperscript{5,6} all four-loop integrals were reduced to a small set of master integrals which were taken from Refs.\textsuperscript{7,8}. Based on these developments a new determination of the quark masses has been performed in Ref.\textsuperscript{9}.

The extraction of $m_Q$ from low moments of the cross section $\sigma(e^+e^- \rightarrow Q\bar{Q})$ exploits its sharp rise close to the threshold for open charm and bottom production and the importance of the contributions from the narrow quarkonium resonances. By evaluating the moments

$$M_r \equiv \int \frac{ds}{s^{n+1}} R_Q(s),$$

with low values of $n$, the long distance contributions are averaged out and $M_r$ involves short distance physics only, with a characteristic scale of order $E_{\text{threshold}}^2 = 2m_Q$. Through dispersion
relations the moments are directly related to derivatives of the vacuum polarization function at $q^2 = 0$,

$$M_n = \frac{12\pi^2}{n!} \left( \frac{d}{dq^2} \right)^n \Pi_Q(q^2) \bigg|_{q^2 = 0},$$  \hspace{1cm} (2)

which can be evaluated in perturbative QCD (pQCD).

The narrow charmonium resonances $J/\Psi$, $\Psi(2S)$ and the higher excitations will obviously contribute to the moments. Open charm production exhibits a sharp rise, nearly like a step function. Beyond the $\Psi(3770)$-resonance a few oscillations are observed which quickly level out into a fairly flat energy dependence. Around and above approximately 5 GeV the cross section is well approximated by pQCD and, furthermore, mass terms can be considered as small corrections\textsuperscript{10,11}. The sensitivity to $m_Q$ is, therefore, concentrated on the small region from $J/\Psi$ up to approximately 5 GeV.

We therefore distinguish three energy regions: First, the region of the narrow resonances $J/\Psi$ and $\Psi(2S)$, second, the "charm threshold region" starting from the D-meson threshold at 3.73 GeV up to approximately 5 GeV, where the cross section exhibits rapid variations and, third, the continuum region where pQCD and local duality are expected to give reliable predictions. For the threshold region we use the data from the BES collaboration\textsuperscript{12,13}, shown in Fig. 1 together with data from MD-1\textsuperscript{14} and CLEO\textsuperscript{15}. Evidently pQCD provides an excellent description of all the data in the continuum region. The description of the perturbative continuum includes the complete mass dependence up to $O(\alpha_s^2)$ plus the dominant mass dependent $O(\alpha_s^3)$ terms\textsuperscript{9} which were used to extrapolate $R_{uds}$ from the region below charm threshold up to 4.8 GeV.

In its domain of analyticity $\Pi_c(q^2)$ can be cast into the form

$$\Pi_c(q^2) = Q^2 \frac{3}{16\pi^2} \sum_{n \geq 0} \bar{C}_n \left( \frac{q^2}{4m_c^2} \right)^n,$$  \hspace{1cm} (3)

where $m_c = m_c(\mu)$ is the \textit{MS} charm quark mass at the scale $\mu$. The perturbative series for the coefficients $C_n$ in order $\alpha_s^2$ was evaluated in Ref.\textsuperscript{16}, the four-loop contributions to $C_0$ and $C_1$ in Refs.\textsuperscript{3,4}. The coefficients depend on $\alpha_s$ and on the charm quark mass through logarithms of the form $ln_m \equiv \ln(m_c^2(\mu)/\mu^2)$. Combining Eqs. (1), (2) and (3), the charm quark mass can be obtained:

$$m_c(\mu) = \frac{1}{2} \left( \frac{C_n}{M_{n^2}} \right)^{1/(2n)},$$  \hspace{1cm} (4)
In the charm threshold region (which includes $\Psi(3770)$) we have to identify the contribution from the charm quark, i.e. we have to subtract the parts arising from the light $u$, $d$ and $s$ quark. In the continuum region above $\sqrt{s} = 4.8$ GeV data are sparse and imprecise. On the other hand, pQCD provides reliable predictions for $R(s)$. Thus in this region we replace data by the theoretical prediction.

We use the results for the moments to obtain in a first step $m_c(3 \text{ GeV})$. The moment with $n = 1$ is least sensitive to non-perturbative contributions from condensates, to the Coulombic higher order effects, the variation of $\mu$ and the parametric $\alpha_s$ dependence. We therefore adopt

$$m_c(3 \text{ GeV}) = 0.986(13) \text{ GeV},$$

as our final result. Transforming this to the scale-invariant mass $m_c(m_c)$ one finds $m_c(m_c) = 1.286(13) \text{ GeV}$. Using the three-loop relation 17,18 between pole- and $\overline{\text{MS}}$-mass this leads to $M_c^{(3\text{-loop})} = 1.666 \text{ GeV}$.

The same approach is also applicable to the determination of $m_b$. Just as in the charm case, a remarkable consistency and stability is observed. For $n = 1$ the error is dominated by the experimental input. For $n = 3$ we obtain $\pm 0.010$ from the experimental input, $\pm 0.014$ from $\alpha_s$ and $\pm 0.006$ from the variation of $\mu$. The three results based on $n = 1, 2$ and 3 are of comparable precision. The relative size of the contributions from the threshold and the continuum region decreases for the moments $n = 2$ and 3. On the other hand, the theory uncertainty estimated from the variation of $\mu$ and the unknown four-loop contribution is still acceptable. Therefore the result from $n = 2$ is taken as the final answer,

$$m_b(10 \text{ GeV}) = 3.609(25) \text{ GeV},$$

corresponding to $m_b(m_b) = 4.164(25) \text{ GeV}$ and a pole mass of $M_b^{(3\text{-loop})} = 4.800 \text{ GeV}$. A comparison with a few selected determinations is shown in Fig. 2.
For various applications, either related to Z-boson decays or in connections to Grand Unified Theories (GUTs), the values of $m_0(\mu)$ at $M_Z$ and $m_0(m_t) = 161.8 \pm 2.0$ GeV (as derived from $M_t = 171.4 \pm 2.1$ GeV\(^{19}\)) are of interest:

$$m_b(M_Z) = 2.834 \pm 0.019 \pm 0.017 \text{ GeV}, \quad m_b(161.8) = 2.703 \pm 0.018 \pm 0.019 \text{ GeV}. \quad (7)$$

The first error reflects the combined error from Eq. (6) and the second one the uncertainty due to $\alpha_s$. The ratio $m_t(m_t)/m_b(m_t) = 59.8 \pm 1.3$ should be a useful input for Grand Unified Theories.

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References

Charm and charmonium spectroscopy at B-factories

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We report on most recent Charm and Charmonium spectroscopy results from the B-factories.

1 Introduction

Since a few years, charm and charmonium spectroscopy has revives, both from experimental and theoretical point of views. Many new states have been discovered triggering numerous theoretical publications. The B-factories with their large enriched charm sample have played a leading role on the experimental side with the observation and study of most of the new states. Other experiments such as CLEO and CDF have also contributed. Classical hadron spectroscopy predicted some of these new states, but not all of them. Therefore a lot of effort have been spent in order to understand the nature of the later. We are summarizing here the most recent and important results in hadron spectroscopy, including strange-charm mesons, charm baryons and charmonium and charmonium-like states.

2 Strange-charm mesons $D_s$

Prior to the B-factories, only 4 strange-strange mesons had been observed: the S-wave states $D_s(1968)^+ (J^P = 0^-)$ and $D_s(2112)^+ (J^P = 0^-)$, and the P-wave states $D_{s1}(2536)^+ (J^P = 1^+)$ and $D_{s2}(2573)^+ (J^P = 2^+)$. In 2003, the CLEO$^1$ and Babar$^2$ collaborations reported two states, the $D_{sJ}(2317)^+$ and the $D_{sJ}(2460)^+$, in continuum production ($e^- e^+ \rightarrow c\bar{c}$). These two states were subsequently observed by Belle$^3$ as well as in B decays. The masses of the states were actually below expectations so there were a lot a speculation whether they were the missing $0^+$ and $1^+$ levels, other ($c\bar{s}$), or exotic states. The $D_{sJ}(2317)^+$ has been observed only in the $D_s^\pm \pi^0$ decay mode, while the $D_{sJ}(2460)^+$ has been seen in the $D_s^{\ast\pm}(D_s^{\ast\pm})\pi^0$, $D_s^{\ast\pm}\gamma$ and $D_s^{\ast\pm}\pi^0\pi^-$ decay modes. Belle$^3$ and Babar$^4$ have performed angular analysis of the $D_{sJ}(2317)^+ \rightarrow D_s^{\ast\pm}\pi^0$ and $D_{sJ}(2460)^+ \rightarrow D_s^{\ast\pm}\gamma$ decays and have shown that the two states are consistent with $J = 0$ and $J = 1$, respectively. A negative parity of the $D_{sJ}(2460)^+$ has been ruled out by an angular analysis of the $D_{sJ}(2460)^+ \rightarrow D_s^{\ast\pm}\pi^0$ decay mode. Searches for isospin partners as well as for doubly charged states and different decay modes$^5$ have been carried out. No signal has been
found. All results indicate that the $\Delta^{*+}(2317)$ and the $\Delta_{sJ}(2460)^+$ are indeed the missing $0^+$ and $1^+$ levels.

In 2006, Babar has reported the observation of a new $D_s^+$ meson\textsuperscript{6} decaying into $D^0 K^+$ and $D^+ K_S^0$ with a mass $m = (2856.6\pm1.5\pm0)\text{MeV}/c^2$ and a width $\Gamma = (48\pm7\pm10)\text{MeV}$, where the first error is statistical and second is systematic. The new state has been observed in two inclusive continuum processes $e^+ e^- \to D^0 K^+ X$ with $D^0 \to K^- \pi^+, K^- \pi^+ \pi^0$ and $e^+ e^- \to D^+ K_S^0 X$ with $D^+ \to K^- \pi^+ \pi^+$. The mass spectra of the three $D$ decay present similar features: a peak at $2.4\text{GeV}/c^2$ due to a reflection of the $D_s(2536)^+$, a clear signal from the $D_s(2573)^+$, a broad structure peaking at a mass of approximately $2700 \text{MeV}/c^2$ and a signal enhancement due to the new state $D_s(2860)^+$. The broad structure at $2700 \text{MeV}/c^2$ is fitted with a mass $m = (2688 \pm 4 \pm 3)\text{MeV}/c^2$ and a width $\Gamma = (112 \pm 7 \pm 36)\text{MeV}$.

This later state was clearly observed in $B^+ \to \bar{D}^0 D_{sJ} \to \bar{D}^0 D^0 K^+$ decay by the Belle collaboration\textsuperscript{7}. It was then assigned to a new $D_{sJ}$ meson, the $D_{sJ}(2700)$. The Dalitz plot analysis shows that the decay $B^+ \to \bar{D}^0 D^0 K^+$ proceeds dominantly through quasi-two-body channels: $B^+ \to \bar{D}^0 D_{sJ}(2700)^+ + r$ and $B^+ \to \psi(3770) K^+$. The measured mass and width are $m = (2715 \pm 11^{+14}_{-11})\text{MeV}/c^2$ and $\Gamma = (115 \pm 20^{+30}_{-20})\text{MeV}$, respectively. The helicity angle distribution favors $J = 1$.

Babar has searched for resonances in $B \to \bar{D}^{(*)} D^{(*)} K$ decays, in 22 decay modes. The $DK$ and $D^*K$ invariant mass distributions have been built summing up all 8 corresponding $B$ decay modes in each case. Both distributions present a clear enhancement near $2700 \text{MeV}/c^2$. However, due to an unknown structure at low mass in the $DK$ invariant mass distribution and to the possible presence of additional resonances in the signal region in the $D^*K$ invariant mass distribution, no attempt has been made to extract the parameters of the the $D_{sJ}(2700)$. A full Dalitz plot analysis is undergoing.

3 Charmonium-like states

The $X(3872)$ was discovered by Belle\textsuperscript{8} in $B^+ \to X(3872) K^+$ with $X(3872) \to J/\psi \pi^+ \pi^-$. It was confirmed\textsuperscript{9} by CDF, D0 and Babar. There have been many speculations about the nature of this state: conventional charmonium state, charmonium hybrid, diquark-antidiquark state\textsuperscript{10}, or $D^0 \bar{D}^{*0}$ molecule\textsuperscript{11}. At present none of the hypothesis is favored. The current mass and width are $m = (3871.2 \pm 0.5)\text{MeV}/c^2$ and $\Gamma < 2.3\text{MeV}$ (90% CL). Babar has found no evidence of a charged partner. Belle\textsuperscript{12} has shown that the $\pi \pi$ invariant mass distribution favors positive parity $P = +1$. The CDF collaboration\textsuperscript{13} has performed an angular analysis of the $X(3872) \to \pi^+ \pi^- \pi^0$ decay and demonstrated that quantum numbers $J^{PC} = 1^{--}$ are favored. In parallel, Belle\textsuperscript{14} and Babar\textsuperscript{15} have observed the decay $X(3872) \to J/\psi \gamma$ which indicates that the charge conjugation number is $C = +1$.

Recently, Belle\textsuperscript{16} has observed an enhancement in the $D^0 \bar{D}^{*0} \pi^0$ system from $B \to D^0 \bar{D}^{*0} \pi^0 K$ decay. The enhancement peaks at a mass $m = (3875.4 \pm 0.7^{+1.2}_{-2.0})\text{MeV}/c^2$. Babar confirmed this observation in the following decays: $B^{+0} \to D^0 D^{*0} K^{+0}$ and $B^{+0} \to D^{*0} D^{0} K^{+0}$ with $D^{*0} \to D^0 \pi^0$ and $D^{*0} \to D^0 K^-$. The ratio of neutral to charged modes branching fractions is $R = 2.23 \pm 0.93 \pm 0.55$. The combined mass from all 4 modes is $m = (3875.6 \pm 0.7^{+1.2}_{-2.0})\text{MeV}/c^2$. Therefore, the mass measurements from Belle and Babar are in very good agreement. However, it is more than 2.5 standard deviations above the mass of the $X(3872)$. Are they the same states or is the state at $3875 \text{MeV}/c^2$ a new resonance?

Babar\textsuperscript{17} has recently discovered a new state, the $Y(4260)$ decaying in to $J/\psi \pi^+ \pi^-$, in the initial state radiation process $e^+ e^- \to \gamma_{ISR} J/\psi \pi^+ \pi^-$. The measured mass and width are $m = (4259 \pm 8 \pm 4)\text{MeV}/c^2$ and $\Gamma = (88 \pm 23 \pm 5)\text{MeV}$, respectively. The quantum numbers are straightforward: $J^{PC} = 1^{--}$. This state is still the subject of many theory papers attempting
to explain its nature: classical charmonium state, tetraquark, or hybrid charmonium. The \( Y(4260) \) was subsequently observed by the Cleo-c, Cleo-III and Belle collaborations. Belle’s mass measurement is higher compared to Babar: \( m = (4295 \pm 10_{-3}^{+11}) \text{MeV}/c^2 \) with \( \Gamma = (133_{-22}^{+13}) \text{MeV} \).

Babar has been searching for \( Y(4260) \rightarrow \psi(2S)\pi^+\pi^- \) decay. The observed \( \psi(2S)\pi \) invariant mass distribution shows an enhancement that is not compatible with the \( Y(4260) \), but with a higher mass resonance. Assuming the enhancement is due to a single resonance, the parameters of this resonance would be \( m = (4324 \pm 24) \text{MeV}/c^2 \) and \( \Gamma = (172 \pm 33) \text{MeV} \) where the errors are statistical only. We shall note that this enhancement is compatible with the \( Y(4260) \) from Belle which measures a higher mass.

Finally, Belle has been reporting three states near \( 3940 \text{MeV}/c^2 \). The \( Z(3930) \) has been observed in \( \gamma\gamma \rightarrow D\bar{D} \), with a mass \( m = (3929 \pm 5 \pm 2) \text{MeV}/c^2 \) and a width \( \Gamma = (29 \pm 10 \pm 2) \text{MeV} \). The angular distribution of the decay strongly favors \( J = 2 \). This state is interpreted as the \((c\bar{c})^2 P_{22}(2^{++}) \) state (the \( \chi''_{c2} \)). The \( Z(3940) \) has been observed in \( B \rightarrow J/\psi\omega(x\pi\pi)K \), with a mass \( m = (3943 \pm 11 \pm 13) \text{MeV}/c^2 \) and a width \( \Gamma = (87 \pm 22 \pm 26) \text{MeV} \). This state has been tentatively assigned to the \((c\bar{c})^2 P_{11}(1^{++}) \) state (the \( \chi''_{c1} \)). The \( X(3940) \) has been observed in continuum production, in the recoil of a \( J/\psi: e^+e^- \rightarrow J/\psi X \), with a mass \( m = (3943 \pm 6 \pm 6) \text{MeV}/c^2 \) and a width \( \Gamma = (15.4 \pm 10.1) \text{MeV} \), and in the \( X \rightarrow D\bar{D}^* \) decay mode. This state is likely to be the \((c\bar{c})^3 S_0(1^{++}) \) state (the \( \eta_{c}(3S) \)).

4 Charm baryon

With the discovery of the \( \bar{\Omega}_{c}^{+30} \), all nine \( J^P = \frac{1}{2}^- \) and six \( J^P = \frac{3}{2}^+ \) ground states (\( L = 0 \)) have now been observed. Several orbitally excited states have already been seen as well.

Babar has observed a new state, the \( \Lambda_{c}(2940)^+ \), decaying into \( D\bar{D}p \), with a mass \( m = (2939.8 \pm 1.3 \pm 1.0) \text{MeV}/c^2 \) and a width \( \Gamma = 17.5 \pm 5.2 \pm 5.9 \text{MeV} \). In this decay mode, Babar confirms the previously reported \( \Lambda_{c}(2880)^+ \) (seen in \( \Lambda_{c}(2880)^+ \rightarrow \Lambda_{c}^{+} \pi^{+}\pi^{-} \)), with a mass \( m = (2881.9 \pm 0.1 \pm 0.5) \text{MeV}/c^2 \) and a width \( \Gamma = (5.8 \pm 1.5 \pm 1.1) \text{MeV} \). The two states have also been reported by Belle, in \( \Lambda_{c}^{+}\pi^{+}\pi^{-} \), with masses and widths \( m = (2881.2 \pm 0.2 \pm 0.4) \text{MeV}/c^2 \) and \( \Gamma = (5.5 \pm 0.5 \pm 0.4) \text{MeV} \), and \( m = (2937.9 \pm 1.0 \pm 1.8) \text{MeV}/c^2 \) and a width \( \Gamma = (10 \pm 5) \text{MeV} \), for the \( \Lambda_{c}(2880)^+ \) and \( \Lambda_{c}(2940) \) respectively. The results from Belle and Babar are in very good agreement. Belle has also performed an angular analysis of the \( \Lambda_{c}(2880)^+ \rightarrow \Sigma_{c}(2455)^{++}/\pi^{0/1} \) and has shown that \( J \geq 5/2 \) is favored.

Belle has observed two new \( \Xi_{c} \) states decaying into \( \Lambda_{c}^{+}\pi^{+}\pi^{-} \), the \( \Xi_{c}(2980)^+ \) with \( m = (2978.5 \pm 2.1 \pm 2.0) \text{MeV}/c^2 \) and \( \Gamma = (43.5 \pm 7.5 \pm 7.0) \text{MeV} \), and the \( \Xi_{c}(3077)^+ \) with \( m = (3076.7 \pm 0.9 \pm 0.5) \text{MeV}/c^2 \) and \( \Gamma = (6.2 \pm 1.2 \pm 0.8) \text{MeV} \). Babar has confirmed these observations and measured the following parameters \( m = (2967.1 \pm 1.9 \pm 1.0) \text{MeV}/c^2 \) and \( \Gamma = (23.6 \pm 2.8 \pm 1.3) \text{MeV} \) and \( m = (3076.41 \pm 0.69 \pm 0.21) \text{MeV}/c^2 \) and \( \Gamma = (6.2 \pm 1.6 \pm 0.5) \text{MeV} \), respectively for the \( \Xi_{c}(2980)^+ \) and the \( \Xi_{c}(3077)^+ \).

The electromagnetic decay \( \Xi_{c} \rightarrow \Xi_{c}\gamma \) has been measured by Babar. It is a confirmation of the Cleo observation. Contributions from \( B \) decays and from continuum are separated with the use of a cut on the \( \Xi_{c} \) momentum measured in the center of mass frame. Branching fractions measured from production in \( B \) decays are \( B(B \rightarrow \Xi_{c}^{+}X) \times B(\Xi_{c}^{+} \rightarrow \Xi^{-}\pi^{+}\pi^{-}) = (1.69 \pm 0.17 \pm 0.10) \times 10^{-4} \) and \( B(B \rightarrow \Xi_{c}^{0}X) \times B(\Xi_{c}^{0} \rightarrow \Xi^{-}\pi^{+}\pi^{-}) = (0.67 \pm 0.07 \pm 0.03) \times 10^{-4} \). For production from the continuum the cross sections are found to be \( \sigma(e^{+}e^{-} \rightarrow \Xi_{c}^{+}X) \times B(\Xi_{c}^{+} \rightarrow \Xi^{-}\pi^{+}\pi^{-}) = 141 \pm 24(exp) \pm 19(model) \text{fb} \) and \( \sigma(e^{+}e^{-} \rightarrow \Xi_{c}^{0}X) \times B(\Xi_{c}^{0} \rightarrow \Xi^{-}\pi^{+}\pi^{-}) = 70 \pm 11(eexp) \pm 6(model) \text{fb} \). The helicity angle distributions of \( \Xi_{c} \) decays are found to be consistent with \( J = \frac{1}{2} \).

Using a large \( \Lambda_{c} \) sample, Babar has studied \( \Lambda_{c}\Lambda_{c} \) correlation production \( e^{+}e^{-} \rightarrow \Lambda_{c}\Lambda_{c}X \), where the \( \Lambda_{c} \) is reconstructed in the \( pK^{-}\pi^{+} \) and \( pK^{*} \) decay modes and the \( \Lambda_{c} \) in the
corresponding charge-conjugate modes. The number of observed events is roughly 4.2 times the number of expected events with respect to models with at least four baryons in the final state. These events show very few additional baryons but multiple mesons, indicating a previously unobserved type of $e^+e^- \rightarrow q\bar{q}$ events.

Conclusion

There have been a lot of activity in charm and charmonium spectroscopy at the B-factories during the last few years. Some of the newly discovered states match theoretical expectations, but most of them, such as the $X(3872)$ and the $Y(4260)$ are still to be understood.

References

New results on two-photon physics from Belle

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Abstract

Results from recent measurements of two-photon reactions in the Belle experiment based on up to 464 fb$^{-1}$ of data are presented. They provide information on various physics aspects related to QCD models of meson and baryon-pair production and properties of charmonia and light-quark resonances.

1 Introduction

Production of hadronic states in two-photon collisions provides valuable information to explore physics of perturbative and non-perturbative QCD and properties of mesons formed by two photon fusion.

At an $e^+e^-$ collider, virtual photons are emitted from the incident electron and positron and collide with each other. The photons are dominated by a component with very small $q^2$, compared with the energy scale of interaction, in case of “zero-tag” measurements where we do not require the detection of the recoil electron/positron and require rather strict transverse-momentum balance of the final-state system (See Sect. 2). We can regard them as compatible to the real photons.

The KEKB is an asymmetric-energy $e^+e^-$ collider whose center-of-mass (c.m.) energy is set at around 10.58 GeV. In the KEKB/Belle experiment$^{1,2}$, we can measure two-photon collisions in the $\gamma\gamma$ c.m. energy ($W$) range 0.8 - 4.5 GeV.

2 Analyses of hadron-pair production processes

We so far measured the cross sections of several processes hadron-pair production in two-photon collisions. The c.m. energy of the incident photon system is derived from the invariant mass of the final hadronic system. We require a rather strict transverse-momentum balance, $|\not{\Sigma}p_{\gamma}| < 0.05 - 0.2$ GeV/c for collecting events only from quasi-real two-photon collisions. This condition is also very effective to get rid of contamination from background processes. We can also study the angular distribution of hadrons in the two-photon c.m. frame.

2.1 $\gamma\gamma \rightarrow K_S^0K_S^0$

We have studied the $\pi^+\pi^-$ and $K^+K^-$ production in two-photon collisions at $W > 2.4$ GeV, previously$^3$. In that work $W$ dependence of the cross sections of the two processes was compared, and the cross section ratio, $\sigma(K^+K^-)/\sigma(\pi^+\pi^-) = 0.89 \pm 0.04 \pm 0.15$ at $W > 3.0$ GeV, was obtained.

We have made a similar analysis for the $K_S^0K_S^0$ production process based on 397.6 fb$^{-1}$ data$^4$. A $K_S^0$ is reconstructed from $\pi^+\pi^-$ requiring a finite decay length under a quite low background conditions. The cut with $|\not{\Sigma}p_{\gamma}|$ is applied at 0.1 GeV/c, and we have obtained the invariant-mass distribution as shown in Fig. 1. It decreases rapidly with the invariant-mass increase, and we find $\chi_{e0}$ and $\chi_{e2}$ charmonia signals with a small background.

We have also measured the angular distributions of kaons in the $\gamma\gamma$ c.m. system. The results are compared with theoretical predictions$^5,6$, and we find that the latter (handbag model) shows a better reproducibility is the angular distributions, although both are consistent with the data.

The cross section of the $K_S^0K_S^0$ process falls much more steeply than in the $\pi^+\pi^-$ or $K^+K^-$ processes with $W$ increase; the preliminary result of the power parameter $n$ in the former processes are $10.6 \pm 0.6 \pm 0.4$, when we fit to $\sigma \sim W^{-n}$ (Fig. 2(a)). In contrast, the
parameter \( n \) is about 7 in the charged-meson processes\(^3\). Accordingly, the cross section ratio \( \sigma(K^0_SK^0_S)/\sigma(K^+K^-) \) is energy dependent in the measurement range and gets very small at the high energies (Fig. 2(b)). The ratio seems to approach the prediction of Benayoun and Chernyak\(^5\) at around 4 GeV.

We derive the products of the two-photon decay width and branching fractions for \( \chi_{c0} \) and \( \chi_{c2} \) to \( K^0_SK^0_S \). They are consistent with our previous measurements of their \( K^+K^- \) decays, considering isospin invariance which predicts the ratio 1:2. These measurements are summarized in Table 1.

### 2.2 \( \gamma\gamma \to \Lambda\bar{\Lambda}/\Sigma^0\bar{\Sigma}^0 \)

Hyperon-pair production processes \( \Lambda\bar{\Lambda} \) and \( \Sigma^0\bar{\Sigma}^0 \) have been studied based on 464 fb\(^{-1}\) data\(^7\). The decay modes \( \Lambda \to p\pi^- \) and \( \Sigma^0 \to \Lambda\gamma \) are used. We select exclusive \( \Lambda\bar{\Lambda} \) candidates with applying a tight \( p_t \)-balance cut (\( |p_t| < 50 \text{ MeV/c} \)) for reducing the background from \( \Sigma^0 \) decays. The exclusive \( \Sigma^0\bar{\Sigma}^0 \) candidates are selected with requiring photons from \( \Sigma^0 \) decays as well as the \( p_t \)-balance in the \( \Sigma^0\bar{\Sigma}^0 \) system.

The invariant-mass distributions for the two processes are shown in Fig. 3. In the \( \Lambda\bar{\Lambda} \) invariant-mass distribution, we find a peak near the \( \eta_c \) mass (\( \sim 2.98 \text{ GeV/c}^2 \)). Another narrow peak near the \( J/\psi \) mass (\( \sim 3.10 \text{ GeV/c}^2 \)) is due to the backgrounds from double initial state radiation events. We find a peak from the \( \eta_c \) also in the \( \Sigma^0\bar{\Sigma}^0 \) invariant-mass distribution.

We summarize the preliminary results of the products of the two-photon decay width of \( \eta_c \) and the branching fractions to the hyperon pairs in Table 1. They are compared with the similar measurement of the \( pp \) mode from our previous publication\(^8\). These measurements can give ratios of the branching fractions of the \( pp \) and the hyperon-pair modes of the \( \eta_c \) decays. They are compatible with unity and may slightly differ from the prediction of a diquark model\(^9\) that gives smaller branching fractions for the hyperon-pair modes.

The energy dependence of the cross sections of the baryon-pair production processes is summarized in Fig. 4. Our preliminary \( \Lambda\Lambda \) and \( \Sigma^0\Sigma^0 \) results are consistent with the L3 result\(^10\) but do not agree well with the CLEO result\(^11\). The cross section of the three processes including the \( pp \) seem to approach each other at \( W > 3 \text{ GeV} \).
2.3 $\gamma\gamma \to f_0(980) \to \pi^+\pi^-$

We study also light-quark mesons. The true nature of $f_0(980)$ is still a mystery; no concrete seat for this particle is prepared in the ordinary $qq$ meson assignment table in spite of its established presence in various processes. It might be an exotic state or be explained by some special theories as relativistic $qq$ models. The two-photon decay width ($\Gamma_{\gamma\gamma}$) should be a crucial key.

We find that the $f_0(980)$ appears as a small peak in the invariant spectrum of $\pi^+\pi^-$ from the two-photon collisions as shown in Fig. 5, where the large $\mu\mu$ background is subtracted with a careful particle identification using an electromagnetic calorimeter.

We fit the $f_0(980)$ resonance line shape and obtain the resonance parameters; $M = 985.6^{+1.2}_{-1.0} +_{-1.6}^{+1.1}$ MeV/$c^2$, $\Gamma_{\pi^+\pi^-} = 34.2^{+13.9}_{-11.8} +_{-2.5}^{+1.2}$ MeV, $\Gamma_{\gamma\gamma} = 205^{+38}_{-33} +_{-17}^{+14}$ eV, for the mass, $\pi^+\pi^-$ and $\gamma\gamma$ partial decay widths, respectively. Although the two-photon decay width is considerably smaller than a prediction based on $d\sigma/dd\Omega$ resonance model, we still cannot conclude whether it is an exotic state. We also make a measurement of angular distributions.

Since we do not find any peak structure near the $\eta'$ mass, we derive an upper limit of the branching fraction of $\eta' \to \pi^+\pi^-$, which is $P$ and $CP$ violating process. New upper limits, $\mathcal{B}(\eta' \to \pi^+\pi^-) < 2.8 \times 10^{-3}$ ($< 3.3 \times 10^{-4}$) have been obtained under the assumption of the interference (no interference) with the non-resonant $\pi^+\pi^-$ process.
Table 1: Products of the two-photon decay width of a charmonium and the branching fraction of its decay from Belle measurements. The first and second errors are statistical and systematic, respectively.

<table>
<thead>
<tr>
<th>Charmonium and decay</th>
<th>$\Gamma_{\gamma\gamma}B$ (eV)</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\chi_{c0}\rightarrow K^0_SK^0_S$</td>
<td>7.80 ± 0.65 ± 0.72</td>
<td>4 prelim.</td>
</tr>
<tr>
<td>$\chi_{c2}\rightarrow K^0_SK^0_S$</td>
<td>0.31 ± 0.05 ± 0.03</td>
<td>4 prelim.</td>
</tr>
<tr>
<td>$\chi_{c0}\rightarrow K^+K^-$</td>
<td>14.3 ± 1.6 ± 2.3</td>
<td>3</td>
</tr>
<tr>
<td>$\chi_{c2}\rightarrow K^+K^-$</td>
<td>0.44 ± 0.11 ± 0.07</td>
<td>3</td>
</tr>
<tr>
<td>$\eta_c\rightarrow \Lambda\Lambda$</td>
<td>6.21 ± 1.01 ± 0.09</td>
<td>7 prelim.</td>
</tr>
<tr>
<td>$\eta_c\rightarrow \Sigma^+\Sigma^0$</td>
<td>9.80 ± 2.50 ± 1.03</td>
<td>7 prelim.</td>
</tr>
<tr>
<td>$\eta_c\rightarrow p\bar{p}$</td>
<td>7.20 ± 1.53 ± 0.17</td>
<td>8</td>
</tr>
<tr>
<td>$Z(3930)\rightarrow DD$</td>
<td>180 ± 50 ± 30</td>
<td>14</td>
</tr>
<tr>
<td>$\chi_{c0}\rightarrow \pi^+\pi^-$</td>
<td>15.1 ± 2.1 ± 2.3</td>
<td>3</td>
</tr>
<tr>
<td>$\chi_{c2}\rightarrow \pi^+\pi^-$</td>
<td>0.76 ± 0.14 ± 0.11</td>
<td>3</td>
</tr>
<tr>
<td>$\chi_{c2}\rightarrow \gamma J/\psi$</td>
<td>114 ± 11 ± 9</td>
<td>15</td>
</tr>
</tbody>
</table>

3 Other results of the $\chi_{cJ}(1P, 2P)$ states

We summarize other results for the $\chi_{cJ}$ charmonium states from previous Belle experiments in Table 1. We have observed a new charmonium state which is tentatively designated as $Z(3930)$ in the $\gamma\gamma \rightarrow DD$ process. We conclude that this is a candidate for the radial-excited state, $\chi_{c2}(2P)$, from the obtained resonance parameters and the angular distribution of the $D$ mesons.

4 Summary

Various two-photon processes in a zero-tag mode have been measured with the Belle detector. The $K^0_SK^0_S$ production process shows much smaller cross section than in the $K^+K^-$ processes above 2.4 GeV in the $\gamma\gamma$ c.m. energy. We also find the $\chi_{c0}$ and $\chi_{c2}$ in the data. $\Lambda$-pair production and $\Sigma^0$-pair production have been measured. The production cross sections are compared with our previous $pp$ results. We also measure the properties of the $\eta_c$ decays to these baryon pairs. We find a cross-section peak from $f_0(980)$ in the $\gamma\gamma \rightarrow \pi^+\pi^-$ spectrum.

4. Belle Collaboration, W. T. Chen et al., hep-ex/0609042 (2006); The results updated recently for publication are found in the v2 (2007).
D Hadronic Analyses at CLEO

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The CLEO-c results on D meson production and hadronic decays obtained with currently available data sets are presented.

1 Introduction

Recent CLEO-c results on D meson production and hadronic decays are presented in this overview. The CLEO-c is a general purpose detector at CESR, Cornell Electron-positron Storage Ring. The detector configuration is an upgraded version of the CLEO III. In order to run at charm production energy the Silicon Vertex Detector was replaced by the 6 stereo-layers drift chamber; magnetic field in the superconducting solenoid is reduced to 1 T for optimal momentum resolution at lower energy. For stable operation at low energy the 12 superconducting wigglers have been installed at CESR upgrade.

The CLEO-c experimental program was started in October 2003 and will continue until April 2008. Analyses included in this overview use three samples of events.
1) At 3770 MeV we have collected luminosity ~560 pb\(^{-1}\), that corresponds to about 4 M produced \(\psi(3770)\) dominantly decaying to \(D\bar{D}\) pairs. The 281 pb\(^{-1}\) of this sample are processed. We plan to accumulate more data at \(\psi(3770)\) by the end of runs.
2) The 12 points scan in the range from 3970 to 4260 MeV with total luminosity of ~60 pb\(^{-1}\) is performed in order to find an optimal energy for \(D_s\) meson study. We find that the \(D^{-}_s D^{*+}\) cross section is maximal around 4170 MeV.
3) At 4170 MeV we have collected 314 pb\(^{-1}\), of which 195 pb\(^{-1}\) are processed. At this energy we plan to accumulate 750 pb\(^{-1}\) in total.

Below I discuss four topics: (i) the absolute \(D^0\), \(D^+\), and \(D_s\) meson hadronic branching fractions measurement; (ii) the \(e^+e^- \rightarrow D^0_{(s)} \overline{D^0_{(s)}}\) cross section measurement in the energy range from 3970 to 4260 MeV; (iii) the inclusive \(\eta, \eta',\) and \(\phi\) meson production branching fractions measurement in the \(D^0\), \(D^+\), and \(D_s\) decays; and (iv) the precision \(D^0\) mass measurement. Non-covered recently published results and ongoing analyses are also listed below.

2 Absolute hadronic branching fractions of \(D^0\) and \(D^+\) mesons

The measurement of absolute hadronic branching fractions of \(D^0\) and \(D^+\) meson decays is early published\(^3\) for luminosity 56 pb\(^{-1}\). Now it is updated for 281 pb\(^{-1}\). We use technique pioneered by MARK III collaboration: count the yield of the single tags (ST), events where the one side
$D$ of $\bar{D}D$ pair is reconstructed in one of considered hadronic modes, the other side $D$ is ignored; count the yield of the double tags (DT), events where both $D$ mesons are reconstructed in considered hadronic modes. Efficiency is defined from full MC simulation. The total number of $\bar{D}D$ pairs produced and branching fractions are extracted from the $\chi^2$ fit. The yields are defined using the signal variables $\Delta E = E_D - E_{\text{beam}}$ and $m_{BC} = \sqrt{E_{\text{beam}}^2 - P_D^2}$, where $E_{\text{beam}}$ is a beam energy, $E_D$ and $P_D$ are the reconstructed $D$ meson candidate energy and momentum, respectively. Typical resolutions are $\sigma(\Delta E) = 7 - 10$ MeV and $\sigma(m_{BC}) = 1.3$ MeV/$c^2$ for modes comprising of tracks only. Presence of $\pi^0$ meson in the final state degrades resolution by a factor of two roughly. The DT yields are extracted from 2D fit to the $m_{BC}(D)$ versus $m_{BC}(D)$ distributions. One of them is shown in Fig. 1 for MC events. Fit accounts for $m_{BC}$ resolutions, beam energy spread, ISR, mis-reconstruction of one-side or both D mesons. Projections of the 2D scatter-plots and relevant yields for $D^0\bar{D}^0$ and $D^0D^-$ events are shown in Fig. 2. The ST yields are extracted from fit to the $m_{BC}$ distributions, Fig. 3. We have considered three decay modes for $D^0$ and six for $D^+$ meson with largest branching fractions. Fit uses ARGUS function for the background shape and “first principles” in order to parameterize the signal component: the $m_{BC}$ resolution, ISR, Breit-Wigner shape for $\psi(3770)$. The total number of reconstructed ST is $\sim 230K$ for $D^0$, and $\sim 167K$ for $D^+$ meson. Results for nine branching fractions are shown in Table 1. We compare them with PDG-2004$^4$ because PDG-2006$^5$ includes our results from 56 pb$^{-1}$ analysis$^3$. The six of measured branching fractions are consistent with the world average values$^4$ within one standard deviation, the three – within two standard deviations. Little has changed from 56 pb$^{-1}$ sample, but systematic uncertainties dominate now. It should be noted that the final state radiation (FSR) effect is included in the efficiency calculation. Without

<table>
<thead>
<tr>
<th>Mode</th>
<th>$B$, %</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^0 \rightarrow K^-\pi^+$</td>
<td>3.87±0.04±0.08</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^-\pi^+\pi^0$</td>
<td>14.6±0.1±0.4</td>
</tr>
<tr>
<td>$D^0 \rightarrow K^-\pi^+\pi^-\pi^+$</td>
<td>8.3±0.1±0.3</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+$</td>
<td>9.2±0.1±0.2</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+\pi^0$</td>
<td>6.0±0.1±0.2</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+\pi^0\pi^+$</td>
<td>1.55±0.02±0.05</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+\pi^-\pi^0$</td>
<td>7.2±0.1±0.3</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+\pi^-\pi^+$</td>
<td>3.13±0.05±0.14</td>
</tr>
<tr>
<td>$D^+ \rightarrow K^-\pi^+\pi^+\pi^-\pi^0\pi^+$</td>
<td>0.93±0.02±0.03</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Mode</th>
<th>$B$, %</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^+<em>s \rightarrow K^{0}</em>{S}K^+$</td>
<td>1.50±0.09±0.05</td>
</tr>
<tr>
<td>$D^+_s \rightarrow K^-K^+\pi^+$</td>
<td>5.57±0.30±0.19</td>
</tr>
<tr>
<td>$D^+_s \rightarrow K^-K^+\pi^+\pi^0$</td>
<td>5.62±0.33±0.51</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \pi^+\pi^-\pi^+$</td>
<td>1.12±0.08±0.05</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \pi^+\pi^-\pi^+$</td>
<td>1.47±0.12±0.14</td>
</tr>
<tr>
<td>$D^+_s \rightarrow \pi^+\eta^+$</td>
<td>4.02±0.27±0.30</td>
</tr>
</tbody>
</table>
account of the FSR the branching fractions decrease up to 2\% (depending on mode).

3 D meson pair production cross sections

In order to find an optimal energy for $D_s$ meson study we performed the energy scan with total luminosity of 60 pb$^{-1}$ in 12 points from 3970 to 4260 MeV. We launched this study because in earlier experiments the total hadronic cross section was only measured. We have measured six pair production cross sections, $D\bar{D}$, $DD^*$, $D^*\bar{D}^*$, $D_s^\pm D_s^{\mp}$, $D_s^{\pm}\bar{D}_s^{\mp}$, and $D_s^{*\pm}D_s^{*\mp}$, shown in Fig. 4. The largest $D_s$ meson production rate can be achieved at energy around 4170 MeV in the process $e^+e^\to D_s^{\pm}\bar{D}_s^{\mp}$, which cross section reaches 0.9 nb.

4 Absolute $D_s$ hadronic branching fractions

Good kinematic separation between the $D$ meson pair combinations can be achieved if we only reconstruct the $D^{0}$ and $D^+$ candidates and ignore $\gamma$ and $\pi^0$ from $D^*$ and $D_s^*$ decays. It is demonstrated in Fig 5, where the $K^+K^-\pi^+$ invariant mass is plotted versus beam constrained mass. Cutting on invariant mass and $m_{BC}$, Fig. 6 (left), we select events of the process $e^+e^-\to D_s^{*}\bar{D}_s^{*}$\mp. In order to measure the $D_s$ meson hadronic branching fractions we use the same
5 Inclusive branching fractions for $\eta$, $\eta'$, and $\phi$

We have measured the branching fractions for the $\eta$, $\eta'$, and $\phi$ mesons inclusive production in $D$-meson decays. These results are published for 281 pb$^{-1}$ at $\psi(3770)$ and 195 pb$^{-1}$ at 4170 MeV. We completely reconstruct the tag-side $D$ meson using three decay modes for $D^0$, five modes for $D^+$, and six modes for $D_s^+$ meson. Using the rest particles on the other side, we inclusively reconstruct the $\eta \to \gamma \gamma$, $\eta' \to \pi^+ \pi^- \eta$, and $\phi \to K^+ K^-$ decays. Invariant masses are used as a signal variables, as shown for example in Fig. 9. Subtracting background we define relevant yields. Measured branching fractions are shown in Table 3. We noticed that the $\phi$ meson production branching fractions depend on $\phi$ momentum. It is also interesting to note how the decay rate depends on presence of the $s\bar{s}$ quarks. The $\phi$ meson is almost pure $s\bar{s}$ state, $\eta'$ has a larger $s\bar{s}$ component than $\eta$. Table 3 reflects the fact that $D_s$ meson prefers to decay in $s\bar{s}$-rich states comparing to the $D^0$ and $D^+$ mesons.
6 \( D^0 \) meson mass measurement

The precision \( D^0 \) mass measurement is recently published\(^7\). Current mass value\(^5\) has 0.4 MeV/\( c^2 \) uncertainty, that is a result of averaging over several experiments, dominated by the LGW, MARK II, and NA32 measurements. In previous experiments the \( D^0 \) mass was measured using the \( D^0 \rightarrow K^- \pi^+ \) and \( D^0 \rightarrow K^- \pi^+ \pi^- \pi^- \) decays. With 281 pb\(^{-1} \) sample we use an advantage of the \( D^0 \rightarrow K^0_S \phi \) decay (mass spectrum is shown in Fig. 7) which has a small energy released, \( M(D^0) - M(\phi) - M(K_S) = 347 \) MeV/\( c^2 \), that restricts a momentum range of the final particles, (400 < \( P_K, P_\pi < 600 \) MeV/\( c \)). We estimate a systematic uncertainties in momentum calibration using inclusively reconstructed \( K^0_S \rightarrow \pi^+ \pi^- \) in \( D \) meson decays. The maximum uncertainty arises due to the momentum and \( \cos \theta \) dependence for \( \pi \)-tracks. The magnetic field calibration is done using invariant mass reconstruction in the decay \( J/\psi \rightarrow \mu^+ \mu^- \). We also check the \( \pi \)-meson momentum calibration using invariant mass reconstruction in the decay \( \psi(2S) \rightarrow \pi^+ \pi^- J/\psi \). In both cases we relay on high precision mass measurement\(^8\) of the \( J/\psi \) and \( \psi(2S) \). We have measured the \( D^0 \) meson mass value, \( M(D^0) = (1864.847 \pm 0.150_{\text{stat.}} \pm 0.095_{\text{syst.}}) \) MeV/\( c^2 \), with statistical and systematic uncertainty a few times smaller than the world average.

7 Non-covered results on \( D \)-meson hadronic decays

Several topics, listed below, have not been covered in this overview. Two papers have been recently published, \(^a\)Branching fraction for the DCSD \( D^+ \rightarrow K^+ \pi^- \pi^\circ \), and \(^a\)Measurement of interfering \( K^{*+}K^- \) and \( K^{*-}K^+ \) amplitudes in the decay \( D^0 \rightarrow K^+K^-\pi^0\pi^- \). The results of
Table 3: Inclusive branching fractions (in %) of D meson decays and their ratios.

<table>
<thead>
<tr>
<th>Mode</th>
<th>$\eta'X$</th>
<th>$\eta X$</th>
<th>$\phi X$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$D^0$</td>
<td>9.5 ± 0.4 ± 0.8</td>
<td>2.48 ± 0.17 ± 0.21</td>
<td>1.05 ± 0.08 ± 0.07</td>
</tr>
<tr>
<td>$D^+$</td>
<td>6.3 ± 0.5 ± 0.5</td>
<td>1.04 ± 0.16 ± 0.09</td>
<td>1.03 ± 0.10 ± 0.07</td>
</tr>
<tr>
<td>$D_s^+$</td>
<td>23.5 ± 3.1 ± 2.0</td>
<td>8.7 ± 1.9 ± 0.8</td>
<td>16.1 ± 1.2 ± 1.1</td>
</tr>
<tr>
<td>$D_s^+/D^0$</td>
<td>2.47 ± 0.34 ± 0.18</td>
<td>3.51 ± 0.80 ± 0.27</td>
<td>15.3 ± 1.6 ± 0.8</td>
</tr>
<tr>
<td>$D_s^+/D^+$</td>
<td>3.73 ± 0.57 ± 0.27</td>
<td>8.37 ± 2.23 ± 0.64</td>
<td>15.6 ± 1.9 ± 0.8</td>
</tr>
</tbody>
</table>

Several analyses are expected to be presented shortly: the $D^0$, $D^+$, and $D_s$ hadronic branching fraction measurements for single and double Cabibbo suppressed decays; the quantum correlation analysis of the $D^0D^0$ decays; the Dalitz plot analyses of the decays $D^+ \to K^-\pi^+\pi^+$, $D^+ \to \pi^-\pi^+\pi^+$, $D_s^+ \to K^-K^+\pi^+$, $D^0 \to K_0^{(S,L)}\pi^+\pi^-$, $D^0 \to K^0_S\pi^0\pi^0$, etc., with and without CP and flavor tags.

8 Summary

In summary, CLEO-c experiment is taking data from October 2003 until April 2008. We have collected 4 M of $D\bar{D}$ pairs, 0.3 M of $D_s^\pm D_s^{\ast\mp}$ pairs at $\sqrt{s} = 4170$ MeV, and 28 M of $\psi(2S)$. We expect to collect more data in the same energy regions by the end of runs at CESR. Preliminary results on $D$ meson production and decays are presented. In most measurements we have reached better precision than the world average.

References

2. CLEO-c and CESR-c A New Frontier of Weak and Strong Interactions, CLNS-01/1742.
III - D0D0bar mixing mini-session

Chairperson: A. Czarnecki
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.) $\times 10^{-3}$ and $y^2 = 9.7 \pm 4.4$ (stat.) $\pm 3.1$ (syst.) $\times 10^{-3}$, 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 CP 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 di-penguin$^2$ diagrams is $O(10^{-8} - 10^{-10})$, 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 $O(10^{-4})$. New physics predictions span the same large range.$^4$ We present evidence for $D$ mixing consistent with these expectations and with previous experimental limits.$^5$ We also compare $D^0$ and $\bar{D}^0$ samples separately, and find no evidence for CP 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$^4\theta_C \approx 0.3\%$ relative to CF decay, where $\theta_C$ is the Cabibbo angle. We distinguish $D^0$ and $\bar{D}^0$ by their production in the decay $D^{*+} \rightarrow \pi^+_D D^0$. In RS decays, the $\pi^+_D$ 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 $\Delta m = m_2 - m_1$ ($\Delta \Gamma = \Gamma_2 - \Gamma_1$) 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

$$
\frac{T_{\text{mix}}}{e^{-\Gamma t}} \propto R_D + \sqrt{|R_D|} y' \Gamma t + \frac{y'^2}{4} (\Gamma s)^2, \tag{1}
$$

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 CF amplitudes. We study both CP-conserving and CP-violating cases. For the CP-conserving case, we fit for the parameters $R_D$, $x^2$, and $y'$. To search for CP violation, we apply Eq. 1 to the $D^0$ and $\overline{D}^0$ samples separately, fitting for the parameters $R_D^\pm$, $x^{2\pm}$, $y'^\pm$ for $D^0(\pm)$ and $\overline{D}^0(-)$ 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 $K^{\pm}\pi^\mp$ invariant mass $1.81 < m_{K\pi} < 1.92$ GeV/c$^2$ and $e^+e^-$ center-of-mass frame (CM) momentum $p_{D0} > 2.5$ GeV/c. We require the $\pi^\pm$ 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 $\pi^\pm$, constraining the $D^0$ daughters to originate from a common vertex while simultaneously requiring the $D^0$ and $\pi^\pm$ 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 $P(\chi^2) > 0.1\%$, $M(K\pi\pi\pi) - m_{K\pi}$ 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 $\pi^\pm$ misreconstructed $D^0$ and combinatorial background. The signal component has a characteristic peak in both $m_{K\pi}$ and $\Delta m$. The random $\pi^\pm$ 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 $\pi^\pm$ 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 $\pi_+^+$ background. Projections of the WS data and fit are shown in Fig. 1.

![Graphs showing $m_{K\pi}$ and $\Delta m$ distributions](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 $\pi_+^+$ 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 $\pi_+^+$ 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 22.9 units. Including the systematic uncertainties, this corresponds to a significance equivalent to 3.9 standard deviations ($1 - \text{CL} = 1 \times 10^{-4}$) 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 $1.843 < m_{K\pi} < 1.883$ 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 $1 - \text{CL} = 0.317$ ($1\sigma$), $4.55 \times 10^{-2}$ ($2\sigma$), $2.70 \times 10^{-3}$ ($3\sigma$), $6.33 \times 10^{-4}$ ($4\sigma$) 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 $R_D = \sqrt{R_D^+ R_D^-}$ and

$$A_D = (R_D^+ - R_D^-)/(R_D^+ + R_D^-)$$

listed in Table 1, from the fitted $R_D^\pm$ values. The best fit points $(x^{2\pm}, y^{2\pm})$ shown in Table 1 are more than three standard deviations away from the no-mixing hypothesis. The shapes of the $(x^{2\pm}, y^{2\pm})$ 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 $\{m_{K\pi}, \Delta m\}$ 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 ± 0.08 ± 0.04</td>
</tr>
<tr>
<td>No CP violation</td>
<td>$x'^2$</td>
<td>3.03 ± 0.16 ± 0.10</td>
</tr>
<tr>
<td></td>
<td>$y'$</td>
<td>-0.22 ± 0.30 ± 0.21</td>
</tr>
<tr>
<td></td>
<td>$y'^+$</td>
<td>9.7 ± 4.4 ± 3.1</td>
</tr>
<tr>
<td>$CP$ violation allowed</td>
<td>$R_D$</td>
<td>3.03 ± 0.16 ± 0.10</td>
</tr>
<tr>
<td></td>
<td>$A_D$</td>
<td>-21 ± 52 ± 15</td>
</tr>
<tr>
<td></td>
<td>$x'^2+$</td>
<td>-0.24 ± 0.43 ± 0.30</td>
</tr>
<tr>
<td></td>
<td>$y'^+$</td>
<td>9.8 ± 6.4 ± 4.5</td>
</tr>
<tr>
<td></td>
<td>$x'^2-$</td>
<td>-0.20 ± 0.41 ± 0.29</td>
</tr>
<tr>
<td></td>
<td>$y'^-$</td>
<td>9.6 ± 6.1 ± 4.3</td>
</tr>
</tbody>
</table>

We validated the fitting procedure on simulated data samples using 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 $m_K$, $\delta m$, proper time, and $\sigma$ PDFs. We varied the $t$ and $\sigma$ requirements. In addition, we considered variations that keep or reject all $D^{**}$ candidates sharing tracks with other candidates. For each source of systematic error, we compute the significance $s^2 = 2 \left[ \ln \mathcal{L}(x'^2, y') - \ln \mathcal{L}(x'^2, y'_i) \right] / 2.3$, where $(x'^2, y')$ are the parameters obtained from the standard fit, $(x'^2, y'_i)$ the parameters from the fit including the $i^{th}$ systematic variation, and $\mathcal{L}$ 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 \(1 + \Sigma x_t^2 = 1.3\) 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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References

EVIDENCE FOR $D^0 - \bar{D}^0$ MIXING

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We observe evidence for $D^0 - \bar{D}^0$ mixing by measuring the difference in apparent lifetime when a $D^0$ meson decays to the $CP$ eigenstates $K^+K^-$ and $\pi^+\pi^-$, and when it decays to the final state $K^-\pi^+$. We find $g_{CP} = (1.31 \pm 0.32\text{(stat.)} \pm 0.25\text{(syst.)})\%$, 3.2 standard deviations from zero. We also search for a $CP$ asymmetry between $D^0$ and $\bar{D}^0$ decays; no evidence for $CP$ violation is found. These results are based on 340 fb$^{-1}$ of data recorded by the Belle detector at the KEKB $e^+e^-$ collider.

1 Introduction

The phenomenon of mixing between a particle and its anti-particle has been observed in several systems of neutral mesons$^{1,2}$: neutral kaons, $K_S^0$, and most recently $B_d^0$ 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$^{3,4}$. The largest predictions are$^4 |x|, |y| \sim \mathcal{O}(10^{-2})$. Loop diagrams including new, as-yet-unobserved particles could significantly affect the experimental values$^5$. $CP$-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$^6$.

Both semileptonic and hadronic $D$ decays have been used to constrain $x$ and $y$$^1$. Here we study the decays to $CP$ eigenstates $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$; treating the decay time
distributions as exponential, we measure the quantity

\[ y_{CP} = \frac{\tau(K^+\pi^-)}{\tau(K^+K^-)} - 1, \]  

(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\(^7\) 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\(^8\); the average value is \( \sim 2 \) standard deviations (\( \sigma \)) above zero. Our measurement yields a nonzero value of \( y_{CP} \) with \( > 3 \sigma \) significance\(^9\). We also search for CPV by measuring the quantity

\[ A_{\Gamma} = \frac{\tau(D^0 \to K^+\pi^-) - \tau(D^0 \to K^-K^+) + \tau(D^0 \to K^-\pi^+)}{\tau(D^0 \to K^-K^+)} \]  

(2)

this observable equals\(^7\) \( A_{\Gamma} = \frac{1}{2}A_M y \cos \phi - x \sin \phi \).

Our results are based on 540 fb\(^{-1}\) of data recorded by the Belle experiment\(^{10}\) at the KEKB asymmetric-energy e\(^+\)e\(^-\) collider\(^{11}\), running at the center-of-mass (CM) energy of the \( \Upsilon(4S) \) resonance and 60 MeV below. To avoid bias, details of the analysis procedure were finalized without consulting quantities sensitive to \( y_{CP} \) and \( A_{\Gamma} \).

2 Event selection

We reconstructed \( D^{*+} \to D^0 \pi^+_s \) decays\(^6\) with a characteristic slow pion \( \pi_s \), and \( D^0 \to K^+K^- \), \( K^-\pi^+ \), 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 \( D^{*+} \) 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 \( D^{*+} \) decay, \( q = (m_{D^0} - M - m_{\pi_s})^2 \). \( M_{D^0} \) is the invariant mass of the \( D^0 \pi_s \) combination and \( m_{\pi_s} \) is the \( \pi_s \) 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 \( \Delta M \equiv M - m_{D^0} \); \( |\Delta q| \equiv q - (m_{D^0} - m_{\pi_s})^2 \); and \( \sigma_q \). 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; we required \( |\Delta M|/\sigma_M < 2.3 \), \( |\Delta q| < 0.80 \) MeV, and \( \sigma_q < 370 \) fs. We found \( 111 \times 10^3 K^+K^- \), \( 122 \times 10^3 K^-\pi^+ \), and \( 49 \times 10^3 \pi^+\pi^- \) signal events, with purities of 98%, 99%, and 92% respectively.

3 Lifetime fit

The relative lifetime difference \( y_{CP} \) was determined from \( D^0 \to K^+K^- \), \( K^-\pi^+ \), and \( \pi^+\pi^- \) decay time distributions by performing a simultaneous binned maximum likelihood fit to the three

\(^a\)Charge conjugate modes are implied unless explicitly stated otherwise.
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,

$$dN/dt = \frac{N_{\text{sig}}}{\tau} \int e^{-(t-t')/\tau} \cdot R(t-t') \, dt' + B(t).$$  \hspace{1cm} (3)

The resolution function $R(t-t')$ was constructed from the normalized distribution of the decay time uncertainties $\sigma_t$ (see Fig. 1). 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 $\sigma_t$, 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. We 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. We therefore choose a parameterization

$$R(t-t') = \sum_{i=1}^{n} f_i \sum_{k=1}^{3} w_k G(t-t'; \sigma_{sk}, t_0),$$  \hspace{1cm} (4)

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. (4). 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.

4 Results

Fits to the $D^0 \rightarrow K^+K^-, K^-\pi^+$ and $\pi^+\pi^-$ data are shown in Fig. 2(a)-(c). The fitted lifetime of $D^0$ mesons in the $K^-\pi^+$ final state, $\tau_{DB} = (408.7 \pm 0.6 \, \text{(stat.)}) \, \text{fs}$, 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

$$y_{CP} = (1.31 \pm 0.32(\text{stat.}) \pm 0.25(\text{syst.})) \%.$$  \hspace{1cm} (5)

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. 2(d), which shows the ratio of decay time distributions for
Figure 2: Results of the simultaneous fit to decay time distributions of (a) $D^0 \to K^+K^-$, (b) $D^0 \to K^-\pi^+$ and (c) $D^0 \to \pi^+\pi^-$ 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 $D^0 \to K^+K^-, \pi^+\pi^-$ and $D^0 \to K^-\pi^+$ decays. The solid line is a fit to the data points.

$D^0 \to K^+K^-, \pi^+\pi^-$ and $D^0 \to K^-\pi^+$ decays. We also searched for CP violation by separately measuring decay times of $D^0$ and $\bar{D}^0$ mesons in CP-even final states. We find an asymmetry consistent with zero, $A_T = (0.01 \pm 0.30\text{(stat.)} \pm 0.15\text{(syst.)})\%$.

References

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QCD and High Energy Hadronic Interactions

IV - Light quark spectroscopy session

Chairperson: A. Czarnecki
INSIGHTS AND PUZZLES IN LIGHT QUARK PHYSICS

H. Leutwyler

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Recent work in light flavour hadron physics is reviewed. In particular, I discuss the significance of the progress achieved with light dynamical quarks on the lattice for the effective low energy theory of QCD. Also, I draw attention to some puzzling results from NA48 and KTeV concerning the scalar form factor relevant for $K_{\mu3}$ decay – taken at face value, these indicate physics beyond the Standard Model.

Keywords: Quantum chromodynamics, light quarks

1 Introduction

At low energies, the most important characteristic of QCD is that the energy gap is very small – a consequence of the fact that the $u$- and $d$-quarks happen to be very light. In the theoretical limit where the quark masses $m_u$ and $m_d$ are set equal to zero, the Hamiltonian of QCD acquires an exact, spontaneously broken symmetry. The spectrum of the theory does then not have an energy gap at all: in the limit, the pions become massless particles, playing the role of the Goldstone bosons which necessarily occur if a continuous symmetry spontaneously breaks down.

The remarkable theoretical progress made in light flavour hadron physics in recent years relies on the fact that these properties can be used to construct an effective theory ("chiral perturbation theory", referred to as $\chi$PT), that allows us to analyze the low energy structure of QCD in a controlled manner. In this framework, all of the Green functions formed with the quark currents can be calculated in terms of the coupling constants occurring in the effective Lagrangian, order by order in the chiral expansion. For recent reviews of $\chi$PT, I refer to Scherer\textsuperscript{1}, Bijnens\textsuperscript{2} and Colangelo\textsuperscript{3} An up-to-date account of our knowledge of the effective coupling constants can be found in a recent conference report by Ecker\textsuperscript{4}.

In my talk, I focused on a few selected aspects of this development. In particular, I emphasized the fact that the progress achieved in the numerical simulation of QCD on a lattice made
it possible to reach sufficiently light quarks, so that the extrapolation to the quark masses of physical interest can be done in a controlled manner, using \( \chi PT \). I will discuss this in some detail below. Another recent development which I will briefly report on, concerns the semileptonic decay \( K \rightarrow \pi \mu \nu \). Both KTeV and NA48 have recently published new results on the scalar form factor of this decay, which are in flat conflict with a venerable low energy theorem, established by Callan and Treiman in the sixties.\(^5\) If these results are confirmed, then the Standard Model is in conflict with observation in one of those reactions which we thought are best understood.

For lack of space, I cannot cover the third topic which I dealt with at these Rencontres: the progress made in understanding the properties of the interaction among the pions. This interaction plays a crucial role in many contexts – the Standard Model prediction for the magnetic moment of the muon is perhaps the most prominent example. To close this introductory section, I briefly list the corresponding keywords and indicate where more information about this can be found. The \( \pi \pi \) scattering amplitude has been calculated to NNL in \( \chi PT \).\(^6\) The resulting representation is very accurate in the interior of the Mandelstam triangle, but in the physical region, the convergence of the series is rather slow. The range of energies where the chiral representation yields meaningful results can be extended with the inverse amplitude trick, but one is then leaving the territory where model independent statements can be made. There is a general method that does not suffer from such shortcomings: in a limited region of the complex plane, dispersion theory imposes a set of exact relations between the real and imaginary parts of the partial wave amplitudes, the Roy equations.\(^7,8\) The region where these equations are valid includes the poles on the second sheet generated by the lowest resonances of QCD: \( \sigma, \rho, \omega, f_0(980), a_0(980) \). The crucial parameters that control the low energy properties of the scattering amplitude in this framework are the two subtraction constants. It is convenient to identify these with the two S-wave scattering lengths, because the low energy theorems of \( \chi PT \) make very sharp predictions for these. Together with the low energy theorems, the Roy equations pin the scattering amplitude down within remarkably small uncertainties.\(^9\) The angular momentum barrier ensures that the S- and P-waves dominate at low energies, but the framework also yields very accurate predictions for partial waves with higher angular momenta. On the basis of this method, it can be demonstrated beyond reasonable doubt that the lowest resonance of QCD carries the quantum numbers of the vacuum and the position of the corresponding pole can be worked out rather accurately.\(^9\) For an overview of these developments, I refer to the proceedings of Chiral Dynamics 2006 and the references quoted therein.\(^10\)

2 Size of the energy gap

In order to illustrate the progress achieved on the lattice, I consider one of the key issues in QCD: understanding the size of the energy gap. The quark masses \( m_u \) and \( m_d \) are very small, but they are different from zero. Accordingly, the Hamiltonian of QCD is not invariant under chiral rotations – the quark mass term breaks chiral symmetry and there is an energy gap: the symmetry breaking equips the Goldstone bosons with a mass. The quark masses \( m_u, m_d \) represent a quantitative measure of the strength of the symmetry breaking. We know that the symmetry breaking is very small, but, unfortunately, the Standard Model does not offer an understanding of why that is so – the entire fermion mass pattern looks bizarre and yet remains to be understood.

As pointed out by Gell-Mann, Oakes and Renner,\(^13\) the square of the pion mass is proportional to the strength of the symmetry breaking, \( M_\pi^2 \propto (m_u + m_d) \). This property can now be checked on the lattice, where – in principle – the quark masses can be varied at will. Fig. 1 shows the result for two recent lattice simulations of QCD with two flavours. In view of the fact that in these calculations, the quarks are treated dynamically, the quality of the data is impressive. The masses are sufficiently light for \( \chi PT \) to allow a meaningful extrapolation to the quark mass
values of physical interest. The results indicate that the ratio $M_\pi^2/(m_u + m_d)$ is nearly constant out to values of $m_u, m_d$ that are about an order of magnitude larger than in nature.

3 Lattice determinations of the effective coupling constants $\tilde{\ell}_3$ and $\tilde{\ell}_4$

The Gell-Mann-Oakes-Renner relation represents the leading term in the expansion in powers of the quark masses. At next-to-leading order, this expansion contains a logarithm:

$$\begin{equation}
M_\pi^2 = M^2 \left(1 + \frac{M^2}{32\pi^2 F^2} \ln \frac{M^2}{\Lambda \tilde{\ell}_3^2} + O(M^4) \right), \quad \tilde{\ell}_3^2 \equiv B(m_u + m_d).
\end{equation}$$

Chiral symmetry fixes the coefficient of the logarithm in terms of the pion decay constant $F$, but does not determine the scale $\Lambda$ of the logarithm. A crude estimate was obtained more than 20 years ago, on the basis of the SU(3) mass formulae for the pseudoscalar octet:

$$0.18 \text{ GeV} < \Lambda_3 < 2 \text{ GeV} \quad \iff \quad \tilde{\ell}_3 \equiv \ln \frac{\Lambda_3^2}{M^2} = 2.9 \pm 2.4.$$  

The logarithmic term implies that the lines in Fig. 1 cannot be straight. For the central value, $\tilde{\ell}_3 = 2.9$, the formula (1) yields only little curvature, but if $\tilde{\ell}_3$ was at the lower (upper) end of the quoted range, the plot of $M_\pi^2$ versus $m_u = m_d$ would be strongly bent upwards (downwards), visibly departing from the lattice results shown in Fig. 1, already at the lowest quark mass value. Evidently, these results strongly constrain the value of the coupling constant $\tilde{\ell}_3$.

| Table 1: Determinations of the effective coupling constants $\tilde{\ell}_3$ and $\tilde{\ell}_4$ |
|---|---|---|---|
| | XPT $^{9,14}$ | MILC $^{17}$ | Del Debbio et al. $^{18}$ | ETM $^{12}$ |
| $\tilde{\ell}_3$ | $2.9 \pm 2.4$ | $0.6 \pm 1.2$ | $3.0 \pm 0.5$ | $3.62 \pm 0.12$ |
| $\tilde{\ell}_4$ | $4.4 \pm 0.2$ | $3.9 \pm 0.5$ | $-$ | $4.52 \pm 0.06$ |

The first row in Table 1 compares the number in equation (2) with recent determinations of $\tilde{\ell}_3$ on the lattice. The second row lists the analogous results for the coupling constant $\tilde{\ell}_4$. 

\[ \text{Fig. 1: Lattice results for } M_\pi^2 \text{ as a function of the quark mass} \]
which controls the quark mass dependence of the pion decay constant. To my knowledge, the first lattice calculation of effective coupling constants based on dynamical quarks was carried out by the MILC collaboration.\textsuperscript{15} In that project, the coupling constants $f_4, L_5, L_6, L_8$, which occur in the effective chiral SU(3)$\times$SU(3) Lagrangian at first nonleading order, were determined by analyzing the quark mass dependence of $M_\pi, M_K, F_\pi$ and $F_K$ by means of $\chi$PT. The corresponding values of $\tilde{\ell}_3$ and $\tilde{\ell}_4$ are readily worked out, using standard one loop formulae.\textsuperscript{16} The numbers quoted in the third column are obtained from a recent update of the MILC analysis, which relies on staggered quarks.\textsuperscript{17} The results obtained by Del Debbio et al.\textsuperscript{18} are based on two flavours of Wilson quarks, while the European Twisted Mass Collaboration\textsuperscript{12} uses two flavours of twisted mass Wilson quarks. All of the numbers are consistent with the values found in $\chi$PT, but some of them are more accurate. Unfortunately, the results obtained with staggered quarks do not agree well with those based on Wilson quarks. Possibly, this is related to the fact that the latter calculations treat all quarks except $u$ and $d$ as infinitely heavy – as emphasized by Stern and collaborators, the strange quark may play a significant role here.\textsuperscript{19} Lattice results for 3 flavours of light Wilson quarks should become available soon, so that it should be possible to identify the origin of the difference.

![Graph showing theoretical and experimental results for the $\pi\pi$ S-wave scattering lengths.](image)

Figure 2: Theoretical and experimental results for the $\pi\pi$ S-wave scattering lengths.

The values of the low energy constants $\tilde{\ell}_3$ and $\tilde{\ell}_4$ play a crucial role also in the theoretical prediction for the $\pi\pi$ scattering lengths. This is illustrated in Fig. 2, where theory is compared with experiment (the lattice results for $\tilde{\ell}_3, \tilde{\ell}_4$ are converted into corresponding values for $a_0^2, a_0^3$ using $\chi$PT and dispersion theory\textsuperscript{9,10}). While the lattice results, DIRAC and the NA48 data on the cusp in $K \to 3\pi$ confirm our predictions for $a_0^2, a_0^3$, the $K_{sl}$ data of NA48/2 give rise to a puzzle: the phase extracted from the transition amplitude deviates from the theoretical prediction for $a_0^2 - a_0^3$. The discrepancy originates in the fact that neutral kaons may first decay into a pair of neutral pions, which then undergoes scattering and winds up as a charged pair. As pointed out by Colangelo, Gasser and Rusetsky, the mass difference between the charged and neutral pions affects this process in a pronounced manner: it pushes the phase of the transition amplitude up by about half a degree – an isospin breaking effect, due almost exclusively to the electromagnetic interaction. The dash-dotted line in Fig. 2, which is taken from the talk of B. Bloch-Devaux at KAOH 2007, shows the likelihood contour ($\chi^2 = \chi_{min}^2 + 2.3$) of the so corrected, preliminary NA48/2 data. The intersection with the region allowed by the low energy theorem for the scalar radius yields $a_0^2 = 0.220(9)$. However, there is a discrepancy with
the E865 data, for which the likelihood contour is shown as a dashed line (kindly provided by G. Colangelo): the isospin correction spoils the good agreement between these data and the prediction. The discrepancy only concerns the region $M_{\pi\pi} > 350$ MeV. While E865 collects all events in this region in a single bin, the resolution of NA48/2 is better. The fit to all $K_{e4}$ data is therefore dominated by NA48/2. For a detailed discussion of these issues, I refer to the talks by B. Bloch-Devaux, G. Colangelo and J. Gasser at KAON 2007. I conclude that the puzzle is gone: $K_{e4}$ confirms the theory to remarkable precision.

4 Puzzling results in $K_{\mu3}$ decay

The low energy theorem of Callan and Treiman\(^5\) predicts the size of the scalar form factor of the decay $K \to \pi \mu \nu$ at one particular value of the momentum transfer, namely $t = M^2_K - M^2_\pi$:

$$f_0(M^2_K - M^2_\pi) = \frac{F_K}{F_\pi} + O(m_u, m_d) .$$

(3)

Within QCD, the relation becomes exact if the quark masses $m_u$ and $m_d$ are set equal to zero. The corrections of first nonleading order, which have been worked out long ago,\(^20\) are tiny; they lower the right hand side by $3.5 \times 10^{-3}$. In the meantime, the chiral perturbation series of $f_0(t)$ has been worked out to NNL.\(^21\) As pointed out by Jamin, Oller and Pich,\(^22\) the curvature of the form factor can be calculated with dispersion theory, so that the prediction for the value at $t = M^2_K - M^2_\pi$ can be converted into a rather accurate prediction for the slope: $\lambda_0 = 0.0157(10)$. Note that, since the $\pi K$ interaction is rather weak, the curvature is small. As witnessed by the very good agreement between the dispersive representations of Jamin et al.\(^22\) and Bernard et al.\(^23\), theory determines the curvature of the form factor, reliably and to good precision.

Very recently, the NA48 collaboration published their analysis of the $K^0_{\mu3}$ form factors.\(^24\) Their result for the scalar slope, $\lambda_0 = 0.0117(7)(1)$, is in flat conflict with the prediction of Jamin, Oller and Pich. Using the parametrization proposed by Bernard et al.\(^23\), NA48 obtains $f_0(M^2_K - M^2_\pi) = 1.155(8)(13) \times f_+(0)$. The value of $F_K/F_\pi$ is sensitive to the theoretical input used for $f_+(0)$. Accurately measured branching ratios and the value of $V_{ud}$ (which has been determined to high precision on the basis of nuclear $\beta$ decays) imply $F_K/F_\pi = 1.242(4) \times f_+(0)$. So, if the NA48 results are correct, then the second term on the right hand side of (3) must amount to a contribution of $0.087(17) \times f_+(0)$. For a correction of $O(m_u, m_d)$, this size is unheard of (as mentioned above, the one loop approximation for this term amounts to -0.0035 and is thus smaller by a factor of about 20). The electromagnetic corrections also need to be investigated, but I very strongly doubt that these could remove the discrepancy.

The NA48 experiment is not the first to measure the slope of the scalar form factor. The first results were obtained in the seventies. In particular, the high statistics experiment of Donaldson et al.\(^25\) had confirmed the theoretical expectations with a slope of $\lambda_0 = 0.019(4)$. More recent experiments, however, came up with quite different results. In particular, three years ago, the KTeV collaboration at Fermilab arrived at a remarkably small scalar slope: $\lambda_0 = 0.0137(2)(131)$. Analyzing 0.54 million charged kaon decays, ISTRA, on the other hand, obtained a much higher value: $\lambda_0 = 0.0183(11)(6)$.\(^27\) If the results of NA48, KTeV as well as those of ISTRA were correct, then the dependence of the form factors on the momentum transfer would have to show very strong isospin breaking — that would be extremely interesting in itself.

My conclusion is that the experimental situation calls for clarification. In particular, an analysis of the charged kaon decays collected by NA48 might help removing the dust. There are not many places where the Standard Model fails. Hints at such failures deserve particular attention.\(^6\)

\(^6\)In the meantime, a preliminary analysis of the KLOE data on the scalar $K^0_{\mu3}$ form factor became available.\(^28\) The result for the slope, $\lambda_0 = 0.0156(26)$, is in excellent agreement with the theoretical prediction, thus removing the puzzle . . . at the price of generating a new one: KLOE is in conflict with NA48.
Acknowledgments

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η PHYSICS AND φ RADIATIVE DECAYS at KLOE


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Here we present KLOE results on the φ meson decays in π^0 π^0 γ, π^+ π^- γ and η π^0 γ, the measurement of the ratio Br(φ → ηγ)/Br(φ → γ) with the estimate of the η' gluonium content and the measurement of the η mass.

1 Introduction

The KLOE experiment is performed at the Frascati φ factory DAΦNE. DAΦNE is a high luminosity e^+,e^- collider working at √s~1020 MeV, corresponding to the φ meson mass. In the whole period of data taking (2001 – 2006) KLOE has collected an integrated luminosity of 2.5 fb^-1, corresponding to about 8 billions of φ produced and 100 millions of η mesons through the electromagnetic decay φ → γγ. The main part of these events are stored on tape, the trigger efficiency ranging from 95% to 100%. The analyses described here are performed on the data collected in the years 2001-2002 corresponding to about 1/5 of all KLOE statistics.
2 \ φ decays to scalars.

2.1 \ φ → f_0γ → π^0π^0γ, \ φ → f_0γ → π^+π^−γ

KLOE has recently published\(^3\) a study of the Dalitz plot of the decay \(\phi → π^0π^0γ\). In order to fit the Dalitz plot we have used two phenomenological models, one based on a Kaon-loop approach\(^3\), KL in the following, and one based on a point-like coupling between the \(φ\) and the scalars \((f_0, a_0)\), “no structure” (NS)\(^7\) in the following. In order to get an acceptable \(χ^2\) value for the fit, the \(σ\) is needed in the KL approach whose parameters we have fixed at \(M_σ = 462\ MeV/c^2\) and \(Γ_σ = 300\ MeV/c^2\), the \(P(χ^2)\) varies from \(10^{-4}\) without \(σ\) and \(14\%\) with \(σ\). Extracting the scalar part of the amplitude from the fitted model we compute \(\text{Br}(φ → f_0γ → π^0π^0γ) = \left[1.07^{+0.04(\text{fit})^{0.000000}}_{-0.04(\text{syst})^{0.000000}}\right] \times 10^{-4}\). 

Isospin symmetry relates the \(π^0π^0\) state to the \(π^+π^-\) state. The background channels in the two final states are instead very different. In both cases there is large interference between the signal and the background, so it is very important to study both final states in order to cross check each other. The study of the \(f_0\) in the \(π^+π^-\) final state has been published in the paper by KLOE\(^3\), in this analysis an earlier version of the KL\(^7\) and the NS model has been used. Both mass values \(m_{f_0}\), and couplings \(R = g_{f_0KK}^2/g_{f_0ππ}^2\) are in good agreement when using the KL fit while they show significant deviation in the NS fit. In general a large fit instability is seen using the NS model. The value of the \(Br\), evaluated as integral of the scalar amplitude, is \(2.0 - 2.4 \times 10^{-4}\), in agreement with the isospin expectation \(Br(f_0 → π^+π^-) \sim 2Br(f_0 → π^0π^0)\). Combining these two values of \(Br's\) we obtain \(Br(φ → f_0γ) = (3.1 - 3.5) \times 10^{-4}\).

2.2 \ φ → a_0γ → π^0π^0γ

The \(π^0π^0γ\) final state has only one interfering background coming from the decay chain \(φ → π^0π^0 → π^0γ\). The contribution of this decay channel is very small due to the small branching ratio \(ρ → γ\). In this case a direct background subtraction is possible in order to extract the \(Br(φ → π^0γ)\). Other large not interfering background is present in the final selection mainly coming from \(e^+e^- → ωπ^0\) with \(ω → π^0γ\), \(φ → f_0γ\) with \(f_0 → π^0π^0\) and \(φ → γγ\) with \(γ → 3π^0\). When 2 photons are lost in the low polar angle region or one or two pairs of photons overlap in a single energy deposit. These background contributions are determined with the use of the MC simulation, after a careful reweighting of each single contribution using background enriched control samples.

The number of events that pass the selection after the background subtraction is \(N_{stg} = 13099 ± 172\), the data analysed correspond to an integrated luminosity of \(L = 413.0 ± 2.5\ \text{pb}^{-1}\). The efficiency of the whole selection chain is \(ε = 37.9\%\), using \(σ_φ = 3900 ± 80\ \text{nb}\) and \(Br(γ → γγ) = 39.38 ± 0.26\%, we obtain:

\[Br(φ → π^0γ) = \frac{N_{stg}}{σ_φ \cdot Br(γ → γγ)} = \frac{6.95 ± 0.09_{\text{stat.}} ± 0.24_{\text{syst.}}}{ε} \times 10^{-4}\]

The systematic error is mainly due to the knowledge of the \(φ\) cross section, the background subtraction and the knowledge of the photon efficiency. In fig.??, a comparison of the value of this branching ratio with several theoretical models is shown, together with the previous measurements.

3 \ η, η' mixing angle.

KLOE has measured\(^7\) the following ratio \(R\) of branching ratios:

\[R = \frac{BR(φ → ηγ)}{BR(φ → ηγ)} = (4.77 ± 0.09_{\text{stat.}} ± 0.19_{\text{syst.}}) \times 10^{-3}\]
Figure 1: Left: Comparison of Br(φ → γη′γ) with theoretical models and previous measurements. Center: ηγ Dalitz plot. Right: m_φ distribution.

Following the reference by Bramon et al, we can write the η/η' wave function as a linear combination of non strange q̅q, strange s̅s quark pairs plus, only for η', \[ |\eta'| = \cos(\varphi_G)\sin(\varphi_F)|q̅q > + \cos(\varphi_G)\cos(\varphi_F)|s̅s > + \sin(\varphi_G)||(q̅q > + s̅s > .\]

The ratio \( R \) can be related to these parameters using the formula:

\[
R = \cos^2(\varphi_F)\cos^2(\varphi_G)\left(1 - \frac{m_z}{m} Z_{NS} \tan \varphi_V \right)^2 \left( \frac{p_{\eta'}}{p_\eta} \right)^3
\]

we fit this value together with the available data on \( \Gamma(\eta' \rightarrow \gamma\gamma)/\Gamma(\pi^0 \rightarrow \gamma\gamma), \Gamma(\eta' \rightarrow \rho\gamma)/\Gamma(\omega \rightarrow \pi^0\gamma) \) and \( \Gamma(\eta' \rightarrow \omega\gamma)/\Gamma(\omega \rightarrow \pi^0\gamma) \) using the theoretical estimates for \( Z_{NS} \) and \( Z_G^2 \). We obtain \( \varphi_F = (39.7 \pm 0.7)^o \) and \( Z_G^2 = \sin^2\varphi_G = 0.14 \pm 0.04 \), with \( P(\chi^2) = 49\% \). Imposing \( \varphi_G = 0 \) the probability \( P(\chi^2) = 1\% \) and the value \( \varphi_F = (39.7 \pm 0.7)^o \).

4 Measurement of the η mass.

The value of the η mass has been recently measured with high precision by two collaborations NA48 (\( m_\eta = 547.843 \pm 0.030 \pm 0.041 \MeV/c^2 \)) and GEMF (\( m_\eta = 547.311 \pm 0.028 \pm 0.032 \MeV/c^2 \)) using different techniques and production reactions. The two measurements differ by more than eight standard deviations from each other. The GEM measurement is in agreement with the older ones while the NA48 measurement is higher. For this reason it is interesting to provide a further measurement of comparable precision in order to clarify the experimental situation.

We measure the mass studying the decay \( \phi \rightarrow \eta\gamma, \eta \rightarrow \gamma\gamma \). A kinematic fit is performed imposing the 4 constraints given by the energy-momentum conservation. Since the photons are just three the fit overconstrains the energies of the photons that are, practically, determined by the position of the clusters in the calorimeter. The inputs of the fit are the energy, the position and the time of the calorimeter clusters, the mean position of the e^+e^- interaction point, the total four-momentum of the colliding e^+e^- . Each of these variables is determined run by run using e^+e^- → e^+e^- events. The absolute \( \sqrt{s} \) scale has been calibrated against the \( m_\phi \) value as measured by CMD-2. The \( \phi \rightarrow \eta\gamma \) events are selected by requiring three energy deposits in the calorimeter and a loose cut on the \( \chi^2 \) of the kinematic fit. The events surviving the cuts are shown in fig. ??, center where the Dalitz plot for \( (m_{\gamma_1\gamma_2, m_{\gamma_1\gamma_2}}) \) is shown. Three bands are clearly visible. The band at low \( m_{\gamma\gamma}^2 \) is given by the \( \phi \rightarrow \pi^+\gamma, \pi^- \rightarrow \gamma\gamma \), while the other two
Table 1: Relative contributions to $m_\eta$ systematic error.

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<tr>
<th></th>
<th>Calorimeter response</th>
<th>Vertex position</th>
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<tbody>
<tr>
<td>Angular stability</td>
<td>26 %</td>
<td>1 %</td>
</tr>
<tr>
<td>Fit bias</td>
<td>70 %</td>
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bands are $\phi \to \eta \gamma, \eta \to \gamma \gamma$ events. With the shown cut in the Dalitz plot we select a pure sample of $\eta, \pi^0 \to \gamma \gamma$ events. The resulting $m_{\gamma\gamma}$ spectrum (fig. ??, right) can be well fitted with a single gaussian with $\sigma \sim 2.1$MeV/c$^2$. In order to evaluate the systematic error we have evaluated the uncertainties on all the variables that are given as input to the fit.

To determine these uncertainties we have used as a control sample the $e^+e^- \to \pi^+\pi^-\gamma$ events, which allows to check the vertex position, the energy response of the calorimeter, the alignment of the calorimeter with the Drift Chamber. A small correction to the value of the mass has been found due to the kinematic fit algorithm. This correction has been evaluated using the MC simulation and half of the correction has been taken as systematic error. Finally the stability of the result respect to different orientations of the 3$\gamma$ plane has been checked, and the variation taken as systematic error.

The various contributions to the systematic error are summarized in table ???. The result is $m_\eta = 547.822 \pm 0.005_{\text{stat}} \pm 0.069_{\text{syst}}$ MeV/c$^2$. In order to check the validity of the whole procedure, the mass of the $\pi^0$ has been measured with the same procedure obtaining $m_{\pi^0} = 134.915 \pm 0.011_{\text{stat}} \pm 0.058_{\text{syst}}$ MeV/c$^2$. This value is in agreement at 1 $\sigma$ level with the PDG value $m_{\pi^0} = 134.9766 \pm 0.0006$ MeV/c$^2$. Our preliminary result is in very good agreement with the NA48 measurement (0.24$\sigma$) and disagrees with GEM by more than 6$\sigma$.

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RECENT RESULTS FROM NA48/2 ON K_{e4} AND K3\pi DECays,
INTERPRETATION IN TERMS OF \pi\pi SCATTERING LENGTHS

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Recent results from the NA48/2 experiment at the CERN/SPS are presented. Large samples of K_{e4} decays have been collected in 2003 and 2004 in the charged (K^+ \rightarrow \pi^+\pi^-e^+\nu) and neutral (K^+ \rightarrow \pi^0\pi^0\nu) modes. Form factors and branching fraction measurements are reported. The \pi\pi scattering lengths \alpha_0 and \alpha_1 can be extracted from the charged K_{e4} phase shift \delta and also from the cusp in the M_{\pi\pi} distribution of K^+ \rightarrow \pi^+\pi^-\pi^0 decays.

1 Introduction

Charged K_{e4} data are of particular interest as they give access to the \pi\pi phase shift \delta = \delta_0 - \delta_1 in absence of any other hadron. The measured variation of the phase shift with the invariant mass M_{\pi\pi} near threshold can be related to the scattering lengths \alpha_0 and \alpha_1 using dispersion relations and data at intermediate energies. More constrained predictions using two-loop Chiral perturbation theory can be compared to the measurement. The neutral K_{e4} mode does not address the same question as the two \pi^0s cannot exhibit a phase difference. However, the measurement of the single form factor \mathcal{F}_0 in the neutral mode can be compared to the one measured in the charged K_{e4} mode. The cusp observed in the M_{\pi^0\pi^0\pi^0} distribution of the K^+ \rightarrow \pi^0\pi^0\pi^+ decays at M_{00} = 2m_{\pi^+} can be explained by \pi^+\pi^- rescattering terms and provides another measurement of the \alpha_0^0 and \alpha_1^0 scattering lengths.

2 Experimental setup

Two simultaneous K^\pm beams were produced by 400 GeV protons from the CERN/SPS, impinging on a beryllium target. The beams were then deflected in a front-end achromat to select momenta in the range (60 \pm 3) GeV/c and focused \sim 200 m downstream at the first spectrometer chamber. A schematic view of the beam line can be found in. The NA48/2 detector main components are: a magnetic spectrometer consisting of a dipole magnet surrounded by two sets
of drift chambers (the momentum of charged decay products is measured with a relative precision of \( \sim 1\% \) for 10 GeV/c tracks) and a 27 radiation length liquid krypton calorimeter used to measure electromagnetic deposits and identify electrons through their \( E/p \) ratio (the energy and position resolutions are \( \sim 1\% \) and \( \sim 1.5 \text{ mm} \) (resp.) for 10 GeV showers).

More than \( 4 \times 10^9 \) \( K^\pm \rightarrow \pi^\pm \pi^- \pi^\pm \) decays, \( \sim 10^8 \) \( K^\pm \rightarrow \pi^0 \pi^0 \pi^\pm \) decays, \( \sim 10^6 \) \( K^\pm \rightarrow \pi^+ \pi^- e^+ e^- \) and \( \sim 4 \times 10^3 \) \( K^\pm \rightarrow \pi^0 \pi^0 e^+ e^- \) decays have been collected in 2003 and 2004. Preliminary results from the analysis of partial statistics are reported here.

3 Charged \( K_{e4} \) decays analysis

From part of the 2003 data, 370000 candidates were selected for three charged tracks topologies, requiring two opposite sign pions, one electron (\( E/p > 0.9 \)) and missing energy and \( p_1 \) (neutrino). The background contamination was estimated from “wrong sign” events (\( \pi^+ \pi^- e^+ e^- \)), which can only be background if one assumes the validity of the \( \Delta S = \Delta Q \) rule.

The background sources are \( K^\pm \rightarrow \pi^+ \pi^- \pi^\pm \) decays with subsequent \( \pi \rightarrow e\nu \) decay or a pion misidentified as an electron and \( K^\pm \rightarrow \pi^\pm \pi^0 (\pi^0) \) decays with subsequent Dalitz decay of a \( \pi^0 \) with an electron misidentified as a pion and photon(s) undetected. The relative level of background to signal is \( \sim 0.5\% \) and has been cross-checked using Monte Carlo simulated events.

The \( K_{e4} \) decay is fully described by the five kinematic Calibb\-\-Maksymowicz variables:\( ^7 \) two invariant masses \( M_{\pi\pi} \) and \( M_{e\nu} \) and three angles \( \theta_{\pi\pi} \), \( \theta_e \) and \( \Phi \) as shown in Figure 1a. Two axial (F,G) and one vector (H) form factors contribute to the transition amplitude and can be expanded in a partial wave expansion of \( s,p,d \) waves.\( ^6 \)

\[ F = F_s e^{i\delta_s} + F_p e^{i\delta_p} \cos \theta_{\pi\pi} + \ldots \quad G = G_p e^{i\delta_p} + \ldots \quad H = H_p e^{i\delta_H} + \ldots \]

Neglecting \( d \) wave terms and assuming the same phase for \( F_p, G_p, H_p \), only one phase \( \delta(q^2) = \delta_s - \delta_p \) and four form factors are left, which are expanded further\( ^7 \) in powers of \( q^2 \): \( M_{\pi\pi}^2 / 4m_e^2 - 1 \):

\[ F_s = (f_s f_s^* q^2 + f_s^* q^2) \quad F_p = (f_p f_p^* + \ldots) \quad G_p = (g_p + g_p^* q^2) \quad H_p = (h_p + \ldots) \]

From the data sample, a total of 150000 equi-populated bins are defined (ten along \( M_{\pi\pi} \), five along \( M_{e\nu} \), five along \( \cos \theta_{\pi\pi} \), five along \( \cos \theta_e \) and twelve along \( \Phi \)). Ten independent four-parameter fits are performed, one for each bin in \( M_{\pi\pi} \). Then a second fit determines the variation of each form factor with \( q^2 \) as shown in Figures 1b and 2.

In the case of the phase \( \delta \), the variation was fitted using a one parameter function, corresponding to the center line of the Universal Band.\( ^2 \) A rather loose constraint between the \( \pi \pi \) scattering lengths \( a_0^0 \) and \( a_0^2 \) is then imposed.

Systematic uncertainties were assessed comparing two independent analyses using different event selections, kaon reconstructions, detector corrections and fit methods. Uncertainties from acceptance, background contamination, electron identification, radiative corrections and possible unknown \( M_{e\nu} \) dependence of the form factors have been investigated. The results are given


Figure 2: Variation of the form factors with $M_{\pi\pi}$ as measured in 10 independent bins. The line is the result of (a): a quadratic fit in $q^2$ for $F_3$, (b): a linear fit in $q^2$ for $G_\rho$ and (c),(d): constant fits for $F_\rho$ and $H_\rho$.

relative to $f_s$, as absolute values can only be obtained through the measurement of the decay rate:

$$
\begin{align*}
\frac{f_3}{f_s} &= 0.169 \pm 0.009_{\text{stat}} \pm 0.034_{\text{syst}} \\
\frac{g_0}{f_s} &= 0.891 \pm 0.019_{\text{stat}} \pm 0.020_{\text{syst}} \\
\frac{f_\rho}{f_s} &= -0.047 \pm 0.007_{\text{stat}} \pm 0.008_{\text{syst}} \\
\frac{h_\rho}{f_s} &= -0.041 \pm 0.027_{\text{stat}} \pm 0.038_{\text{syst}} \\
\end{align*}
$$

A comparison with $a_{\ell0}^0$ values obtained by previous $K_{e4}$ experiments and using the same Universal band center line constrain shows consistency within the measurement errors.

4 Neutral $K_{e4}$ decays analysis

From 2003 (2004) data, ~10000 (30000) events were selected requiring one charged track (identified as electron), two neutral pions (reconstructed from four $\gamma$'s in the calorimeter) and missing energy and $p_T$ (neutrino). The background contamination was estimated from data by reversing some of the selection requirements. It is mainly due to $K^\pm \rightarrow \pi^0 \pi^0 \pi^\pm$ events where the $\pi^\pm$ is misidentified as an electron. Another contribution comes from radiative $K_{e3}$ decays and an accidental extra photon, the two photons faking a $\pi^0$. The relative background to signal level was 3% in 2003 and 2% in 2004. The reconstructed kaon mass is shown in Figure 3a.

The branching fraction is measured from 2003 data only. The preliminary value, normalized to $K^\pm \rightarrow \pi^0 \pi^0 \pi^\pm$ is: $BR(K_{e4}^{00}) = (2.587 \pm 0.026_{\text{stat}} \pm 0.019_{\text{syst}} \pm 0.029_{\text{ext}}) \times 10^{-5}$, where the external error comes from the uncertainty on the normalization mode branching ratio. As there are two identical $\pi^0$s in the decay, one form factor describes the decay in the s-wave. This form factor and its $q^2$ dependence have been measured from the full statistics: $f_3/f_s = 0.129 \pm 0.036_{\text{stat}} \pm 0.020_{\text{syst}}$ and $f_\rho^0/f_s = -0.040 \pm 0.034_{\text{stat}} \pm 0.020_{\text{syst}}$. While measured with a smaller precision than in the charged $K_{e4}$ mode, the values are in good agreement.

5 Cusp in $K^\pm \rightarrow \pi^0 \pi^0 \pi^\pm$ decays

Preliminary results using part of the data statistics have been already published. However, some improvements in the analysis deserve to be mentioned. The invariant $\pi^0 \pi^0$ mass squared distribution, $M^2_{\pi^0}$, shows a sudden change of slope at $M^2_{\pi^0} = 2m^2_{\pi^0}$ as seen in Figure 3b. This effect has been explained and computed with a rescattering model using the Dalitz plot.
formulation: $M_{2 \phi} = 1 + \frac{1}{2} g_{\phi \phi} + \frac{1}{2} h' \nu^2 + \frac{1}{2} k' \nu^2$ where $k'$ was set to zero, as suggested by\(^{12}\). A fit of the single mass distribution allows the simultaneous determination of $g_{\phi}$, $h'$, $a_{\phi}^2$ and $(a_{\phi}^2 - a_{\phi}^2)$. Going to a 2D-fit requires to use both squared invariant masses $M_{2 \phi}^2$ and $M_{2 \phi}^2$. An alternative choice is $M_{2 \phi}^2$ and $\cos \theta$, where $\theta$ is the angle between the charged $\pi^\pm$ and the $\pi^0$'s direction in their rest frame. The following results were obtained, assuming a zero value for $k'$:

$g_{\phi} = 0.045 \pm 0.004_{\text{stat}} \pm 0.006_{\text{syst}}, \quad h' = -0.047 \pm 0.012_{\text{stat}} \pm 0.013_{\text{syst}},$

$a_{\phi}^2 = -0.041 \pm 0.022_{\text{stat}} \pm 0.014_{\text{syst}}, \quad a_{\phi}^2 - a_{\phi}^2 = 0.266 \pm 0.010_{\text{stat}} \pm 0.008_{\text{syst}} \pm 0.013_{\text{syst}}.$

The systematic uncertainty includes acceptance and trigger efficiency. The external error corresponds to an estimate of the effect of the missing higher order terms and radiative corrections in the rescattering model. When fitting also the $k'$ parameter in the 2D-fit, it is measured significantly away from zero: $k' = 0.0097 \pm 0.0003_{\text{stat}} \pm 0.0008_{\text{syst}}$. The corresponding changes for $g_{\phi}$ and $h'$ are $\sim 2\%$ (resp. $2.5\%$) but negligible for the scattering lengths.

6 Conclusion

Using part of the data recorded in 2003 and 2004, NA48/2 has improved measurements of the $K_{e4}$ form factors in the charged and neutral modes. The neutral $K_{e4}$ branching fraction has been measured. First evidence for a non-zero $\nu^2$ term in the $K^\pm \rightarrow \pi^0 \pi^0 \pi^\pm$ Dalitz plot has been reported. A preliminary value of the $\pi \pi$ scattering length $a_{\phi}^2$ has been obtained with $3\%$ statistical and $3\%$ systematic precision in the conservative approach of the Universal Band. A joint study of the $K_{e4}$ and cusp analyses will provide stringent constraints in the $(a_\phi^2, a_\phi^2)$ plane.

References

III - *Light quark spectroscopy session*

Chairperson: A. Czarnecki

D0D0bar mixing mini-session
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 SLAC. We find the mixing parameters $x^2 = -0.22 \pm 0.30$ (stat.) $\pm 0.21$ (syst.) $\times 10^{-3}$ and $y^2 = 0.7 \pm 0.4$ (stat.) $\pm 0.1$ (syst.) $\times 10^{-3}$, 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 $CP$ 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 di-penguin$^2$ diagrams is $O(10^{-8} - 10^{-10})$, 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 $O(10^{-4})$. New physics predictions span the same large range.$^4$ We present evidence for $D$ mixing consistent with these expectations and with previous experimental limits.$^5$ We also compare $D^0$ and $\bar{D}^0$ samples separately, and find no evidence for $CP$ 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 $\tan^4 \theta_C \approx 0.3\%$ relative to CF decay, where $\theta_C$ is the Cabibbo angle. We distinguish $D^0$ and $\bar{D}^0$ by their production in the decay $D^{*+} \rightarrow \pi^+_D D^0$. In RS decays, the $\pi^+_D$ 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 $\Delta m = m_2 - m_1$ ($\Delta \Gamma = \Gamma_2 - \Gamma_1$) 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 \cite{12}

$$\frac{T_{wa}}{e^{-\Gamma t}} \propto R_D + \sqrt{R_D} y' \Gamma t + \frac{x'^2 + y'^2}{4} (\Gamma t)^2,$$

(1)

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 CF amplitudes. We study both CP-conserving and CP-violating cases. For the CP-conserving case, we fit for the parameters $R_D, x'^2$, 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 $R_D \pm, x'^2, y'$ for $D^0(+) \bar{D}^0(-)$ 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 \cite{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 $p_{CM} > 2.5$ GeV/c. We require the $\pi^+_s$ 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 $\pi^+_s$, constraining the $D^0$ daughters to originate from a common vertex while simultaneously requiring the $D^0$ and $\pi^+_s$ 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 \cite{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 $\pi^+_s$, misreconstructed $D^0$, and combinatorial background. The signal component has a characteristic peak in both $m_{K\pi}$ and $\Delta m$. The random $\pi^+_s$ 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 $D^0 \to K^-\pi^+$ 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 $\pi^+_s$ event class shapes shared between RS and WS datasets.
We find 1,141,500 ± 1,200 RS signal events and 4,030 ± 90 WS signal events. The dominant background component is the random π⁺π⁻ 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/c², and b) $\Delta m$ for WS candidates with $1.843 < m_{K\pi} < 1.883$ GeV/c². 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 π⁺π⁻ 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 π⁺π⁻ 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 ($-2\Delta \ln L$) 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 $-2\Delta \ln L$ 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 $-2\Delta \ln L$ for no-mixing is 21.9 units. Including the systematic uncertainties, this corresponds to a significance equivalent to 3.9 standard deviations $(1 - \text{CL} - 1 \times 10^{-4})$ 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 $\overline{D}^0$ WS candidates and fit result integrated over the signal region $1.843 < m_{K\pi} < 1.883 \text{ GeV/}c^2$ and $0.1445 < \Delta m < 0.1465 \text{ GeV/}c^2$. 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$^{-3}$ (3$\sigma$), 6.33 $\times$ 10$^{-4}$ (4$\sigma$) and 5.73 $\times$ 10$^{-7}$ (5$\sigma$), calculated from the change in the value of $-2\ln L$ 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 $R_D = \sqrt{R_D^2}^{-1}$ and $A_D = (R_D^1 - R_D)/(R_D^1 + R_D)$ listed in Table 1, from the fitted $R_D^1$ values. The best fit points $(x'^{2}\pm y'^{\pm})$ shown in Table 1 are more than three standard deviations away from the no-mixing hypothesis. The shapes of the $(x'^{2}\pm, y'^{\pm})$ CL contours are similar to those shown in Fig. 3. All cross-checks indicate that the close agreement between the separate $D^0$ and $\overline{D}^0$ fit results is coincidental.

As a cross-check of the mixing signal, we perform independent $\{m_{K\pi}, \Delta m\}$ 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 ± 0.08 ± 0.04</td>
</tr>
<tr>
<td>No CP violation</td>
<td>$x^2$</td>
<td>3.03 ± 0.16 ± 0.10</td>
</tr>
<tr>
<td></td>
<td>$y'$</td>
<td>-0.22 ± 0.30 ± 0.21</td>
</tr>
<tr>
<td>CP violation allowed</td>
<td>$R_D$</td>
<td>3.03 ± 0.16 ± 0.10</td>
</tr>
<tr>
<td></td>
<td>$A_D$</td>
<td>-21 ± 52 ± 15</td>
</tr>
<tr>
<td></td>
<td>$x^{2+}$</td>
<td>-0.24 ± 0.43 ± 0.30</td>
</tr>
<tr>
<td></td>
<td>$y'^+$</td>
<td>9.8 ± 6.4 ± 4.5</td>
</tr>
<tr>
<td></td>
<td>$x^{2-}$</td>
<td>-0.20 ± 0.41 ± 0.29</td>
</tr>
<tr>
<td></td>
<td>$y'^-$</td>
<td>9.6 ± 6.1 ± 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.

![Graph](image)

Figure 4: The WS branching fractions from independent $m_K^*$, $\Delta m$ 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 1.5; for the no-mixing hypothesis (a constant WS rate), the $\chi^2$ is 24.0.

We validated the fitting procedure on simulated data samples using 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 $m_{K^*}$, $\Delta m$, proper time, and $\sigma_1$ PDFs. We varied the $t$ and $\sigma_1$ requirements. In addition, we considered variations that keep or reject all $D^+$ candidates sharing tracks with other candidates. For each source of systematic error, we compute the significance $s^2_i = 2 \left[ \ln \mathcal{L}(x^{2i}, y^i) - \ln \mathcal{L}(x^{2i}, y'_i) \right] / 2.3$, where $(x^{2i}, y^i)$ are the parameters obtained from the standard fit, $(x^{2i}, y'_i)$ the parameters from the fit including the $i^{th}$ systematic variation, and $\mathcal{L}$ the likelihood of the standard fit. The factor 2.3 is the 68% confidence level for 2 degrees of free-
To estimate the significance of our results in \( (x^2, y') \), we reduce \(-2\Delta \ln \mathcal{L}\) by a factor of \(1 + \Sigma x_i^2 = 1.3\) 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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References

EVIDENCE FOR $D^0 - \bar{D}^0$ MIXING

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We observe evidence for $D^0 - \bar{D}^0$ mixing by measuring the difference in apparent lifetime when a $D^0$ meson decays to the CP eigenstates $K^+ K^-$ and $\pi^+ \pi^-$, and when it decays to the final state $K^- \pi^+$. We find $g_{CP} = (1.31 \pm 0.32 \text{(stat.)} \pm 0.25 \text{(syst.)})\%$, 3.2 standard deviations from zero. We also search for a CP asymmetry between $D^0$ and $\bar{D}^0$ decays; no evidence for CP violation is found. These results are based on 540 fb$^{-1}$ of data recorded by the Belle detector at the KEKB $e^+ e^-$ collider.

1 Introduction

The phenomenon of mixing between a particle and its anti-particle has been observed in several systems of neutral mesons$^{1,2}$: 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$^{3,4}$. The largest predictions are$^4 |x|, |y| \sim \mathcal{O}(10^{-2})$. Loop diagrams including new, as-yet-unobserved particles could significantly affect the experimental values$^5$. 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$^6$.

Both semileptonic and hadronic $D$ decays have been used to constrain $x$ and $y$.$^1$ Here we study the decays to CP eigenstates $D^0 \rightarrow K^+ K^-$ and $D^0 \rightarrow \pi^+ \pi^-$; treating the decay time
distributions as exponential, we measure the quantity

\[ y_{CP} = \frac{\tau(K^-\pi^+)}{\tau(K^+K^-)} - 1, \]  

(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. Our measurement yields a nonzero value of \( y_{CP} \) with \( > 3\sigma \) significance.\(^9\) We also search for CPV by measuring the quantity

\[ A_\Gamma = \frac{\tau(D^0 \to K^-K^+) - \tau(D^0 \to K^+K^-)}{\tau(D^0 \to K^-K^+) + \tau(D^0 \to K^+K^-)}, \]  

(2)

this observable equals \( A_\Gamma = \frac{1}{2} A_M y \cos \phi - x \sin \phi \).

Our results are based on 540 fb\(^{-1} \) of data recorded by the Belle experiment\(^10\) at the KEKB asymmetric-energy \( e^+e^- \) collider,\(^11\) running at the center-of-mass (CM) energy of the \( Y(4S) \) resonance and 60 MeV below. To avoid bias, details of the analysis procedure were finalized without consulting quantities sensitive to \( y_{CP} \) and \( A_\Gamma \).

2 Event selection

We reconstructed \( D^{*+} \to D^0 \pi_\pm \) decays\(^8\) with a characteristic slow pion \( \pi_\pm \), and \( D^0 \to K^+K^- \), \( K^-\pi^+ \), 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 \( \pi_\pm \) track to originate from the \( e^+e^- \) interaction region. A \( D^{*+} \) momentum greater than 2.5 GeV/\( c \) (in the CM) was required to reject \( D^0 \) 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 \( D^{*+} \) decay, \( q = (M_{D^+} - M - m_\pi)^2 \). \( M_{D^+} \) is the invariant mass of the \( D^0 \pi_\pm \) 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 \( \Delta M \equiv M - m_{D^0} \); \( |\Delta q| \equiv q - (m_{D^+} - m_{D^0} - m_\pi)^2 \); 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: we required \( |\Delta M|/\sigma_M < 2.3 \), \( |\Delta q| < 0.80 \) MeV, and \( \sigma_t < 370 \) fs. We found \( 111 \times 10^3 \) \( K^+K^- \), \( 122 \times 10^3 \) \( K^-\pi^+ \), and \( 49 \times 10^3 \) \( \pi^+\pi^- \) signal events, with purities of 98\%, 99\%, and 92\% respectively.

3 Lifetime fit

The relative lifetime difference \( y_{CP} \) was determined from \( D^0 \to K^+K^- \), \( K^-\pi^+ \), and \( \pi^+\pi^- \) decay time distributions by performing a simultaneous binned maximum likelihood fit to the three

\(^{a}\)Charge conjugate modes are implied unless explicitly stated otherwise.
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,

$$dN/dt = \frac{N_{\text{sig}}}{\tau} \int e^{-t'/\tau} \cdot R(t - t') \, dt' + B(t).$$

The resolution function $R(t - t')$ was constructed from the normalized distribution of the decay time uncertainties $\sigma_t$ (see Fig. 1). 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 $(t_{\text{rec}} - t_{\text{gen}})/\sigma_i$ (where $t_{\text{rec}}$ and $t_{\text{gen}}$ are reconstructed and generated MC decay times), is not well-described by a Gaussian. We 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. We therefore choose a parameterization

$$R(t - t') = \sum_{i=1}^{n} \sum_{k=1}^{3} f_i w_k G(t - t'; \sigma_{ik}, t_0),$$

with $\sigma_{ik} = s_k \sigma_k^{\text{pull}} \sigma_i$, 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. (4). 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.

4 Results

Fits to the $D^0 \rightarrow K^+K^-, K^-\pi^+$ and $\pi^+\pi^-$ data are shown in Fig. 2(a)-(c). The fitted lifetime of $D^0$ mesons in the $K^-\pi^+$ final state, $\tau_{DB} = (408.7 \pm 0.6 \, \text{(stat.)}) \, \text{fs}$, 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

$$\gamma_{CP} = (1.31 \pm 0.32(\text{stat.}) \pm 0.25(\text{syst.})) \%. \tag{5}$$

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. 2(d), which shows the ratio of decay time distributions for
Figure 2: 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.

$D^0 \rightarrow K^+K^-, \pi^+\pi^-$ and $D^0 \rightarrow K^-\pi^+$ decays. We also searched for CP violation by separately measuring decay times of $D^0$ and $\bar{D}^0$ mesons in CP-even final states. We find an asymmetry consistent with zero, $A_T = (0.01 \pm 0.30\,\text{(stat.)} \pm 0.15\,\text{(syst.)})\%$.

References

V - QCD Session

Chairpersons: K. Sliwa and N. Glover
QCD on-shell recurrence relations and the space-cone gauge

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Abstract

We first give a field theoretical derivation of the tree level on-shell gluon (BCFW) recursion relations, by reassembling the tree level gluon Feynman diagrams in a convenient gauge, space-cone. The significance of these recurrence relations is that they allow obtaining multi-gluon tree level amplitudes from the previously computed lower n-point functions. Our proof of the BCFW recursion relations hinges on an algebraic identity in momentum space which is the Fourier transform of Veltman's largest time equation. Then we show that the use of the space-cone gauge is instrumental in selecting the right analytic continuation of the one-loop gluon amplitudes, namely that analytic continuation which leads to an extension of the BCFW recursion relations to the one-loop level.

1 Introduction

QCD calculations are notoriously tedious if one is to follow the usual Feynman-Dyson expansion in some commonly used, such as Feynman, gauge. Over the past few years, great strides have been made to simplify such endeavors. The results for the complete amplitudes at the tree or one-loop level can be quite compact.

Following Witten's proposal for a description of perturbative Yang-Mills gauge theory as a string theory on twistor space [1], and subsequent proposal for an alternative to the usual Feynman diagrams in terms of the so-called maximally helicity violating (MHV) vertices [2], a new set of methods was available for the computation of QCD amplitudes. The latest advance in the form of recursion relations [3, 4], in conjunction with the attendant rules for their construction, is particularly appealing. It is quite obvious from the flavor of such an approach that it bears on the cutting rules in field theory. In fact, some work at the one-loop level under the heading of cut-constructibility clearly points to the same origin [5].

As is well-known, unitarity of the S-matrix and the feasibility of an ordering of a sequence of space-time points are intimately related. Indeed, the ordering need not be with respect to time, as is conventionally done. All that is essential in a perturbation series is that one must be able to separate the positive frequency and the negative frequency components in a propagator according to the signature of a certain linear combination of components $\Delta z$ of the four vector between the two space-time points. For our purpose, a component $\eta \cdot \Delta z$ of the light-cone variables will be a convenient start, where $\eta$ is a light-like vector. We shall rely on the existence of tubes of analyticity to continue such variables into the space cone, in order to incorporate a gauge condition for QCD. The resulting ordering is the equivalent of the largest time equations. The outcome, for QCD in particular, is that one factorizes a physical amplitude into products of physical amplitudes, with some momenta shifted but still on-shell. This is the content of the BCFW recursion relations [3, 4]:

$$A(P, \{P_i\}, Q, \{Q_i\}) = \sum_{i, j} A_L(P, \{P_i\}, K) \frac{1}{(P + \sum_i P_i)^2} A_R(K, \bar{Q}, \{Q_i\}),$$

where $A_L, A_R$ are lower n-point functions obtained by isolating two reference gluons with shifted momenta, $\bar{P} = P - z\eta$, $\bar{Q} = Q + z\eta$ with $\eta^2 = \eta \cdot P = \eta \cdot Q = 0$, on the two sides of the cut. The shifting is necessary in order to preserve energy-momentum conservation. We would like to take this opportunity to point out that in so far as factorization is concerned, the masses of the internal propagators have no bearing. However, the demand that the shifted momenta, which will be called reference momenta, should be on-shell will force these external momenta to be light-like.

We now turn to the important step of gauge fixing. In order to facilitate a natural cancellation of terms at every level of a QCD calculation, the gauge that is most convenient for us is the space-cone gauge [6]. A crucial advantage of this gauge is that when we shift the momenta to obtain recursion relations, the dependence on momenta of the vertices will not be affected. Thereupon the factorization of the amplitudes is the same as that in a scalar theory. It is this special attribute which makes the program manageable.
The BCFW recurrence relations were extended to on-shell one-loop level recurrence relations by Bern, Dixon and Kosower [7]. The BCFW proof relied on a certain analytic continuation of the on-shell amplitude in the complex plane, followed by the evaluation of the integral \( \oint A(z) dz/z \). The latter step localizes the tree level amplitude onto the residues \( Res A(z) \). Since the poles of the tree level amplitudes are simple, and since at the tree level the analytic continuation of BCFW is such that \( A(z) \to 0 \) as \( z \to \infty \), it follows that \( Res A(z) \) correspond to setting in sequential order various internal lines on-shell, thus leading to the on-shell recurrence relations. The one-loop story proved to be more complicated. Since the loop amplitudes have cuts, the choice of the integration contour must be such that it is a deformation of the contour at infinity into a contour which hugs closely each cut. Therefore, only the rational parts of the one loop multi-gluon amplitudes will be determined recursively. Another complication comes from the choice of analytic continuation. In certain one-loop amplitudes, the BCFW analytic continuation prescription generates a boundary term, which is a consequence of having \( A(z) \to \infty \) as \( z \to \infty \). However the problem of selecting the right analytic continuation prescription (without knowing the amplitude beforehand) is greatly simplified by the use of the space-cone gauge. We simply require that the analytic continuation, which is essentially a shift of the gluon external momenta into the complex plane, is such that the space-cone vertices are left unchanged. Since the internal propagators necessarily acquire a \( z \)-dependence, it is not difficult to see that the boundary term is eliminated.

2 On-shell tree level recursion relations from Feynman diagrams

According to the spinor helicity formalism, null vectors can be decomposed into a product of two commuting spinors (twisters): \( P^{\alpha \beta} = p^\alpha p^\beta \equiv |p| \bar{|p|} \). Moreover, we can use the twisters to define a basis on the space of four-vectors

\[
P = p^+ |+||+| + p^- |￢||+| + p^- |￢||+| (1)
\]

As shown by Chalmers and Siegel, [6], in the spinor helicity formalism, a powerful simplification is achieved in the Feynman diagrams by choosing the space-cone gauge: \( a = 0 \), followed by the elimination of the "auxiliary" component \( \phi \) from its equation of motion. The gauge fixed Lagrangian has now only two scalar degrees of freedom

\[
\mathcal{L} = Tr \left[ \frac{1}{2} a^+ \Box a^- - i \left( \frac{\partial}{\partial a^+} \right) [a^+, \bar{\partial} a^-] - i \left( \frac{\partial}{\partial a^-} \right) [a^-, \bar{\partial} a^+] + [a^+, \bar{\partial} a^-] \frac{1}{\Box} [a^-, \bar{\partial} a^-] \right]. \tag{2}
\]

Choosing the space-cone gauge amounts to selecting two of the external momenta to be the reference null vectors for defining a twistor basis: \( |+| |+| |￢| |￢| \), such that the space-cone gauge fixing is equivalent to \( N \cdot A = 0 \), where the null vector \( N \) is equal to \( |+| |￢| \).

2.1 The causality ("largest time") equations

As stated in the Introduction, the recursion relations are rooted in the largest time equation. To this end, we briefly revisit here the causality equations as derived by Veltman [8], but appropriately rewriting them in a light-cone frame.

First, we introduce the following set of rules:
- duplicate the Feynman diagram \( 2^N \) times, for \( N \) vertices, by adding circles around vertices in all possible ways;
- each vertex can be circled or not; a circled vertex will bring a factor of \( i \), and an uncircled vertex will bring a factor of \( (-i) \);
- the propagator between two uncircled vertices is \( \Delta(x-y) \), while the propagator between two circled vertices is the complex conjugate \( \Delta^*(x-y) \);
- the propagator between a circled \( x_k \) and an uncircled \( x_l \) is \( \Delta^+(x_k-x_l) \), while the propagator between an uncircled \( x_k \) and a circled \( x_l \) is \( \Delta^-(x_k-x_l) \).

Clearly, the uncircled Feynman diagram is the usual one, while the fully circled diagram corresponds to its complex conjugate.

The largest time equation states that the sum of all \( 2^N \) circled Feynman diagram vanishes:

\[
F(x_i) + F^*(x_i) + F(x_i) = 0, \tag{3}
\]

where \( F(x_i) \) stands for the usual Feynman diagram, \( F^*(x_i) \) is its complex conjugate, and \( F(x_i) \) is the sum of \( 2^N - 2 \) diagrams in which at least one vertex is circled and at least one is uncircled.
Other causality equations can be obtained by singling out 2 vertices, $x_k$ and $x_l$:

$$F(x_k) = -F(k, l, x_k) - \theta((x_l - x_k)^\alpha)F(k, l, x_l) - \theta((x_k - x_l)^\alpha)F(k, l, x_k),$$  \hspace{1cm} (4)

where $F(k, l, x_i)$ is the sum of all diagrams with neither $k, l$ circled, but at least one other vertex circled, $F(k, l, x_i)$ is the sum of all amplitudes with $k$ uncircled, but $l$ circled and finally, $F(k, l, x_i)$ has $k$ circled and $l$ uncircled.

The $\eta$-shift of two of the external momenta, required in the on-shell recursion relations, is a consequence of the momentum inflow associated with the step-functions which arise in the largest time equation

$$\theta((x_l - x_k)^\alpha)^+ = \int \frac{dz}{2\pi i(z - i\epsilon)} e^{i\pi\eta(x_l - x_k)}. \hspace{1cm} (5)$$

2.2 Reassembling Feynman diagrams into BCFW recursion relations

A crucial observation is that by performing the BCFW shifts of the external momenta of the two space-cone gauge reference gluons, the momentum dependence of the vertices in any Feynman diagram is left unchanged. Then, after factoring out the vertices, the Feynman diagrammatic proof of the tree level on-shell recursion relations is based on an algebraic identity involving the shifted propagators.

To exploit the full generality of the problem, we derive the identity satisfied by the momentum space scalar propagators working under the assumption that we deal with massive propagators, with arbitrary masses.

We single out two external momenta which do not land on the same vertex as reference vectors and call them $P_a$ and $P_b$. For a tree graph, there is a unique path through some of the internal lines which connects $P_a$ to $P_b$. The factorization procedure is to cut these $Q_i$ successively by shifting them by $z\eta$. The on-shell conditions $Q_i^2 + m_i^2 = 0$, $Q_i \equiv q_i - z\eta$ will give us a set of solutions, points in the complex plane, namely $z_i = \frac{Q_i^2 + m_i^2}{2z_i Q_i}$.

More precisely stated, the factorization amounts to splicing the graph into a sum of products of two on-shell graphs with shifted momenta $\{P_a - z_i\eta, \ldots, Q_1\}$ and $\{-Q_1\ldots, P_b + z_i\eta\}$, where $\ldots$ stand for the other momenta in the left graph segment and similarly for those in the right graph segment, with the propagator $\frac{1}{Q_i^2 + m_i^2}$ as the partition.

Imposing the condition that the shifted momenta remain on-shell requires that $P_a \cdot \eta = P_b \cdot \eta = 0$, in other words, $P_a, P_b$ must be null.

The identity which we want to establish is

$$\frac{1}{Q_1^2 + m_1^2 Q_2^2 + m_2^2} \cdots \frac{1}{Q_{n-1}^2 + m_{n-1}^2} = \frac{1}{Q_1^2 + m_1^2 (Q_2 - z_1\eta)^2 + m_2^2} \cdots \frac{1}{(Q_{n-1} - z_1\eta)^2 + m_{n-1}^2}$$

$$+ \frac{1}{(Q_1 - z_2\eta)^2 + m_1^2 Q_2^2 + m_2^2} \cdots \frac{1}{(Q_{n-1} - z_2\eta)^2 + m_{n-1}^2} + \cdots$$

$$+ \frac{1}{(Q_1 - z_{n-1}\eta)^2 + m_1^2} \cdots \frac{1}{(Q_{n-2} - z_{n-1}\eta)^2 + m_{n-2}^2 Q_{n-1}^2 + m_{n-1}^2},$$  \hspace{1cm} (6)

and its proof relies on the partial fractioning formula $\int \frac{dz}{z(x-z_1)(x-z_2)\cdots(x-z_n)} = 0$. Next we notice that eqn. (6) (with massless propagators) is precisely the identity needed to reassemble a generic tree level gluon Feynman diagram into lower on-shell amplitudes.

The connection with the largest time equation comes through the observation that by Fourier transforming (4), we shift the momenta of two gluons landing on the vertices $x_k$ and $x_l$, as a consequence of the step functions in (4). The graph splicing is induced by the $\Delta^\pm$ propagators.

3 On-shell loop recurrence for all same helicity gluons

We begin by making the observation that in the space-cone gauge a one loop same helicity gluon amplitude is built only out of 3-point vertices. To be specific, let us consider 1-loop amplitudes with positive helicity external gluons. Then all the vertices will be trivalent: $\{+, +, -\}$.

First, we notice that we cannot specify the space-cone vector $\eta$ = $|+\rangle |\rangle$ completely in terms of external gluon momenta. The reason is that the external (on-shell) gluons have all the same helicity. We do have an alternative, though: we can select three external gluons and shift their momenta according to

$$P_1 = |1\rangle |1\rangle \rightarrow \tilde{P}_1 = |1\rangle |1\rangle - (|1\rangle + z|2\rangle |1\rangle) |1\rangle$$

$$P_2 = |2\rangle |2\rangle \rightarrow \tilde{P}_2 = |2\rangle |2\rangle - (|2\rangle + z|3\rangle |2\rangle) |2\rangle$$

$$P_3 = |3\rangle |3\rangle \rightarrow \tilde{P}_3 = |3\rangle |3\rangle - (|3\rangle + z|12\rangle |3\rangle) |3\rangle$$  \hspace{1cm} (7)
The sum of the three momenta does not change under the holomorphic shift, because of the Schouten identity: $[33][1] + [31][2] + [12][3] = 0$. Incidentally, we notice that this is the same holomorphic shift which Risager [13] used to prove the CSW rules starting from the BCFW recursion relations. It is also clear, by inspection, that the shifted momenta remain on shell. Lastly, since the vertices for amplitudes with all external gluons of the same helicity are all of the type $(+ - +)$, then these shifts will not change the vertices, nor the external line factors.

Armed with this observation we can proceed to derive the on-shell recurrence relation following BCFW and compute $\oint dz A(z)$, where $A(z)$ is the amplitude evaluated with the shifted momenta. We know that the all plus one-loop amplitude is a rational function, since it vanishes at the tree level. We have also seen that the vertices do not change with the shift (7), but the internal propagators do, with at least one of them being affected. Then we infer that $A(z) \sim 1/z^n$, with $n \geq 1$ as $z \to \infty$. Therefore, $\oint dz A(z) = 0$ when the contour is taken at infinity. In other words, there is no boundary term to contend with. This means that we can factorize the all plus amplitude into lower $n$-point functions associated with the residues of $A(z)$:

$$A_{n}^{(1)}(P_1, P_2, P_3, P_4, \ldots P_n) = A_{3}^{(0)}(\hat{P}_1, \hat{P}_2, K) \frac{1}{2P_1 \cdot P_2} A_{n-1}^{(1)}(K, \hat{P}_3, P_4 \ldots P_n)$$

$$+ A_{3}^{(0)}(\hat{P}_2, \hat{P}_3, K) \frac{1}{2P_2 \cdot P_3} A_{n-1}^{(1)}(K, P_4 \ldots P_n, \hat{P}_1)$$

$$+ A_{3}^{(0)}(\hat{P}_3, P_4, K) \frac{1}{2P_3 \cdot P_4} A_{n-1}^{(1)}(K, P_5 \ldots P_n, \hat{P}_1, \hat{P}_2)$$

$$+ A_{3}^{(0)}(P_n, \hat{P}_1, K) \frac{1}{2P_n \cdot P_1} A_{n-1}^{(1)}(K, \hat{P}_2, P_3, P_4 \ldots P_{n-1})$$

(8)

where the superscripts 0, 1 indicate whether the one-shell amplitude is tree or one-loop level, the boldface letters denote the triplet of external gluon momenta, and the hats denote the shifts made such that the line cut is put on-shell. In each of the four terms $K$ is placed on-shell by the appropriate $z$ shift.

References


HEAVY FLAVOUR PRODUCTION AT HERA

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Studies of charm and beauty production in $ep$ collisions with a center-of-mass energy of 318 GeV are reported from the two HERA collaborations, H1 and ZEUS. The analyses make use of both the HERA-I data sample recorded between 1996 and 2000 and a sample from HERA-II, which started in 2003. The cross sections measured by both H1 and ZEUS experiments are compared with next-to-leading order QCD calculations. The measurement of the charm and beauty contributions to the proton structure function is also presented. The comparison to next-to-next-to-leading order (NNLO) calculations shows agreement within the errors.

1 Introduction

Heavy quark production processes provide a powerful insight into the understanding of Quantum Chromodynamics. The large mass of the heavy quark makes the perturbative calculations reliable, even for total cross sections, by cutting off infrared singularities and by setting a large scale at which the strong coupling can be evaluated. At HERA heavy flavour production is possible both in photoproduction and deep inelastic scattering (DIS) reactions, the latter having dramatically smaller cross sections. While in photoproduction the photon virtuality $Q^2$ is very small ($Q^2 \sim 0$), and the photon is almost real, in DIS $Q^2$ can reach values much higher than the squared quark mass $m_q^2$. In direct-photoproduction processes the quasi-real photon enters directly in the hard interaction whilst in resolved-photoproduction processes the photon acts as a source of partons that take part in the hard interaction.

The dominant process for heavy quark production in DIS and in direct photoproduction in $ep$-collisions at HERA is the boson-gluon fusion (BGF) mechanism, $\gamma g \rightarrow Q\bar{Q}$. At leading order (LO), the BGF process is directly sensitive to the gluon content of the proton. In resolved photoproduction it is necessary to also consider quark excitation diagrams, $Qg \rightarrow Qg$, where the heavy quarks originate from the photon, and the gluon-gluon fusion, $gg \rightarrow Q\bar{Q}$. In photoproduction both the direct and the resolved components contribute to heavy quark production. Calculation tools are available up to next-to-leading order ($\alpha_s^2\text{NLO}$) in the form of Monte Carlo integration programs $^{1,2}$. They use the massive scheme $^3$ in which $u$, $d$ and $s$ are the only active flavours in the proton, and charm and beauty are dynamically produced in the hard scatter. In another (massless) approach, beauty and charm are treated as massless and included in the PDFs. The HERA measurements shown here are compared to massive and massless NLO QCD predictions.
2 Charm Production

The most recent results by the H1 collaboration regarding the $D^{*\pm}$ photoproduction\textsuperscript{4} used a data sample five times larger than in the previous publications\textsuperscript{5}. The $D^*$ meson was detected via the decay channel\textsuperscript{6} $D^{*\pm} \rightarrow D^{\pm}_1 \pi^\pm \rightarrow K^\pm \pi^\pm \pi^\pm$. Details of the heavy quark production process were investigated by studying events with a jet not containing the $D^*$ meson ($D^* + jet$). In Fig. 1(left) the measured differential cross section as a function of the transverse momentum of the $D^*$, $p_T(D^*)$, is compared to NLO calculations based on the massive scheme\textsuperscript{6,7} (FMNR) and a general-mass variable-flavour-number scheme\textsuperscript{8,9} (GMVFNS). The cross section falls steeply with increasing $p_T(D^*)$ as predicted by all calculations. In Fig. 1(right) the differential cross section as a function of the difference in the azimuthal angle between the $D^*$ and the other jet $\Delta \phi(D^*, jet)$ is shown. The results are compared with NLO prediction from FMNR and with predictions based on the zero-mass variable-flavour-number scheme\textsuperscript{10,11} (ZMVFN). A large fraction of the produced $D^* + jet$ combinations deviates from a back-to-back configuration. The NLO calculations are not able to describe the small $\Delta \phi$ behaviour of the data, indicating the presence of higher order contributions in this particular region.

Recent results on $D^*$ production in DIS by the ZEUS Collaboration\textsuperscript{12} are shown in Fig. 2(left). The differential cross section in $Q^2$ is compared to the previous HERA-I result\textsuperscript{13}. Good agreement is seen for the entire $Q^2$ range over which the differential cross section falls by about four orders of magnitude. The cross section is reasonably well described by the NLO calculation which use the ZEUS NLO QCD fit. In (Fig. 2)(left) also shown are the ZEUS results on the $D^*$ cross section in the range $0.05 < Q^2 < 0.7$ GeV$^2$\textsuperscript{14}. The beam-pipe calorimeter of ZEUS was used for the measurement of the scattered lepton, which allows the first measurement of the transition region between photoproduction and DIS. The NLO calculations describe well also this region.

3 Beauty Measurements

H1 has recently measured charm and beauty cross sections using a fit to the lifetime signature of charged particles in jets\textsuperscript{15}. This inclusive method yields measurements of differential cross sections that extend to larger values of transverse momenta than in previous HERA analyses, in which leptons from beauty quark decays were used to measure beauty cross sections. Fig. 3 shows the measured cross sections as a function of the transverse momentum, $p_T^{jet}$, of the leading

\textsuperscript{a}Charge conjugate states are implicitly implied.
Rencontres de Moriond 2007

Figure 2: (left) Differential $D^*$ cross section as a function of $Q^2$ compared to the NLO calculation of HVQDIS. The HERA-II data (solid points) are shown compared to the most recently published ZEUS measurements (open squares). The solid line gives the predictions from the ZEUS NLO QCD fit for $m_c = 1.35$ GeV with the shaded band indicating the uncertainty in the prediction. The results using the BPC data are also reported (open circles). On the right the summary of the latest beauty cross section measurements using different tagging techniques by ZEUS and H1 is shown.

jet. Taking into account the theoretical uncertainties, the beauty cross sections are consistent both in normalisation and shape with a perturbative QCD calculation to next-to-leading order. In Fig. 2(right) a summary of the beauty cross section measurements from both H1 and ZEUS collaborations using different tagging methods is reported. At low transverse momentum values, the data tends to be slightly above the NLO QCD predictions. In this region the cross sections are extracted using double tagging techniques which allow to lower the kinematic threshold due to lower background. HERA II data will be needed to improve the cross section determination in the low- and high-$p_T$ region.

4 $F_2^c$ and $F_2^b$

Measurements of the charm and beauty contributions to the inclusive structure function $F_2$ have been performed recently at HERA. The measurement of $F_2^c$ and $F_2^b$ has been done in a kinematic region where the extrapolation needed to correct for the full phase space is as small as possible. In Fig. 4(left) a summary of the $F_2^c$ measurements as a function of $Q^2$ for different $x$ values is shown. The measurement of $F_2^b$ has been performed by ZEUS for the first time. In Fig. 4(right) $F_2^b$ measured by the two experiments are compared with theoretical predictions, based on fixed-flavour and variable-flavour number schemes; a first NNLO calculation is also reported in the figure. The measurements from the two experiments are compatible within the errors and in agreement with the theory.
Figure 4: (left) $F_2^{ee}$ and (right) $F_2^{gg}$ measurements as a function of $Q^2$ for different $z$ values.

5 Summary

Recent results on beauty and charm production in $ep$ collisions have been presented. The NLO QCD predictions describe well the charm data in a large range of $Q^2$, including the transition region between PHP and DIS. The beauty data agree with the NLO predictions at high-$p_T$ whilst at low-$p_T$ there is a tendency of the data to be slightly above the central NLO predictions. The latest measurements of the $F_2^{ee}$ and $F_2^{gg}$ have been reported.

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Beauty and Charm Production Measurements at the Tevatron

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Results for the production of charm and beauty quarks in p\bar{p} collisions at $\sqrt{s} = 1.96$ TeV (Tevatron) with the two multi-purpose experiments CDF and D0 using an integrated luminosity of up to 1 fb$^{-1}$ are presented. With the data measurement of the production mechanism for charm and beauty quarks are done. These measurements are leading to a better understanding of the QCD in the transition region between perturbative and nonperturbative QCD. The charm-charm angular correlation is measured with D mesons at the CDF experiment. The Quarkonium production for charm and beauty final state at D0 and CDF is discussed.

1 Introduction

The production of long lived heavy quarks, the charm (c) and beauty (b) quark, in hadron-hadron collisions is an active field of research in Quantum Chromo Dynamics (QCD). In the theoretical treatment the heavy quark mass provides a scale just at the transition between non-perturbative and perturbative QCD. The measurement of cross sections, angular quark-quark correlations and polarisation improves our understanding of the QCD transition region. This leads to an excellent program for testing one of the fundamental forces.

Charm and beauty quarks are produced in huge numbers in proton anti-proton collision at a center-of-mass energy of $\sqrt{s} = 1.96$ GeV at the Tevatron collider at Fermilab. CDF$^1$ and D0$^2$ are the two multi purpose detectors at Fermilab. The two experiments have partly complementary utilities for measurements of heavy quark final states. The large angular coverage for muons and the high muon trigger bandwidth of the D0 detector lead to an excellent efficient selection specially for di-muon hadronic final states like the Upsilon.

The measurements of the charm and beauty cross sections were made with small fraction of the available data. For the charm cross section the CDF experiment showed that using a luminosity of about 6 pb$^{-1}$ the measured value to be about a factor of two above the theoretical expectation. This is still well inside the scale uncertainty of the theory$^3$. Also new and updated measurements by both collaborations with data using a luminosity up to 1 fb$^{-1}$ for the beauty cross
Figure 1: The $D^0D^{*-}$ (left) and $D^+D^{*-}$ (right) pair cross section measured by the CDF experiment as a function of $\Delta \phi$. The measurements are compared to Pythia (black line). Also the different contributions from quark pair production (red), flavour excitation (green) and gluon splitting (blue) are shown.

sections show the same tendency: the measured cross sections are above the central values of the improved theoretical expectation but well covered inside the theoretical uncertainties$^{4,5}$.

2 Charm correlation measurements

In QCD, the flavour conservation implies, that charm quarks are always produced as quark anti-quark pairs. The azimuthal angle ($\phi$) correlation of the mesons carrying the two heavy quarks in a event gives a possibility to study the underlying production mechanism in detail$^{6}$. Prompt heavy quark pair production leads to back-to-back production. Heavy quarks produced by splitting of massless gluons have a small $\Delta \phi$. The third production process, the production by flavour excitation leads to a big separation in the rapidity $\eta$ and most of the time one of the two quarks is produced in the forward region.

The CDF collaboration measured the angle correlation in D-Meson$^{7}$ production with an integrated luminosity of 1 fb$^{-1}$. The angular correlation of more than 2000 events in the two signal modes $D^0D^{*-}$ and $D^+D^{*-}$ was investigated. The $D^0$ and $D^+$ are used for the trigger selection of charm events. The slow pion from the $D^{*-}$ decay leads to a high purity of the samples. Simulation studies have shown that the event selection efficiency factorize into the two single efficiencies of the charm Mesons. Therefore the measured single charm cross sections can be used for the overall normalisation. Figure 1 shows the angular correlation in the $D^0D^{*-}$ and in the $D^+D^{*-}$ channel and compares it to the Pythia simulation. While Pythia gives a fair estimate for the overall production, it leads to an underestimation (overestimation) of the contribution from pair production (gluon splitting).

3 Quarkonium

Prior the first measurements at Fermilab, the production of Quarkonium was described by the color singlet model$^8$. However the first measurements at Fermilab showed that this model underestimates the direct production cross section by an order of magnitude for the $J/\Psi$ and by a factor of 50 for the $\Psi(2S)$ production. To address these problems, the color octet model was introduced$^9$. Adjustable hadronisation parameters in this model allow to resonable describe the amplitude and $p_T$-dependence of the production. Recent, different approaches using Pomeronic ideas and $k_T$-factorisation models were introduced to describe the hadronization problem$^{10,11}$. Both models predict that at sufficient high $p_T$ the $J/\Psi$ and the $\Psi(2S)$ will have a longitudinal
The D0 experiment selected the $\Upsilon(1S), \Upsilon(2S)$ and $\Upsilon(3S)$ via di-muon decay and measured the $\Upsilon(1S)$ differential cross section\textsuperscript{12}. Figure 2(left) shows the di-muon mass spectra for two different bins in the pseudorapidity $\eta$ of the final state. Clear significant signals of $\Upsilon(1S)$ and $\Upsilon(2S)$ are observed. Figure 2 (right) shows the measured differential cross section of the $\Upsilon(1S)$. Because of the high $\eta$ acceptance of the D0 muon system these measurement is done in three $\eta$ bins. No significant $\eta$ dependence is observed. It is planned to extend this analysis to a polarisation measurement using a data set with a luminosity of 1 fb$^{-1}$. To extract the polarisation of all three resonances simultaneously an unfolding method is necessary.

The CDF collaboration measured the polarisation for the two vector mesons $J/\Psi$ and $\Psi(2S)$\textsuperscript{13,14} using an integrated luminosity of about 800 pb$^{-1}$. Both modes were selected using the decay into two muons. The flight direction of the $\mu^+$ relative to the flight direction of the vector meson in the proton-proton rest frame of the $p_T$ measured by the polar angle $\theta^*$ which depends on the polarisation parameter $\alpha$ were $\alpha = 1(-1)$ for transversal (longitudinal) polarisation. Figure 3 shows the results for prompt $J/\Psi$ (left) and the $\Psi(2S)$ (right) production. For $J/\Psi$ production the measured polarisation is longitudinal and significantly increases with the $p_T$. This agrees fully with the expectation of the new models. Also the $\Psi(2S)$ polarisation shows a trend towards longitudinal polarisation with higher $p_T$.

Additionally the CDF collaboration measured the relative cross section of the $\chi_1$ and $\chi_2$ mesons\textsuperscript{15}. These mesons decay into $J/\Psi$ mesons via the radiation of a low energetic photon. The precision of the measurement of $\sigma(\chi_2)/\sigma(\chi_1) = 0.70 \pm 0.04$ (stat.) $\pm 0.04$ (sys.) $\pm (0.06$ branching fractions) for the prompt production sets a new standard for this measurement. It excludes a naive estimation based on counting of the different spin orientations which leads to an expected cross-section ratio of 5:3.
4 Summary

The study of heavy flavour production mechanism is an active research topic at the Tevatron collider. Both experiments CDF and D0 have now collected a large sample of data with a luminosity of about 2 fb\(^{-1}\). With this huge data sample it will be possible to access the details of the production mechanism by the measurement of polarisation of quarkonium or the quark anti-quark correlation. This additional information is especially important because the theoretical knowledge of the cross section are dominated by scale uncertainties. The new measurements provide an additional input information for the understanding of the production mechanism for open and hidden quark production in the transition region between perturbative and non-perturbative QCD.

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Production of photons and heavy quarks in CDF

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Abstract

Production of photons in association with heavy quarks at hadron colliders is interesting per se as well as in the light of looking for deviations due to new physics. As is often the case in hadron colliders, the trigger plays a very important role in the measurement. Two strategies can be followed to perform this measurement: either using an unbiased photon trigger, where the only triggered object is the photon, necessarily with a high threshold, or combining photon and jet information. In CDF, this latter approach is possible thanks to the impact parameter trigger that allows enriching the heavy flavor content of a sample triggering on tracks with large impact parameter. Results of measurements performed using both unbiased and impact parameter trigger will be presented.

1 Introduction

The photon + b system is very interesting since its production cross section is very dependent on the b Pdf’s, and is expected to be very sensitive to the presence of new physics due to its b final state. As an example, it is expected to receive contributions from SUSY models with production of chargino-neutralino pairs $\chi_1^{-} \chi_1^{0}$ decaying into $b\gamma \gamma \chi_1^{0}$ through intermediate stop production. From the experimental point of view, this measurement requires identification of b-jets using secondary vertex tag, and of photons from the large $Z^{0}$ background. None of these can be performed with a sufficient purity to allow for an event-by-event selection, therefore statistical separation based on fitting relevant distributions has to be performed.

2 Trigger and analysis strategy

This analysis has been performed using two complementary trigger strategies: the first is to select events based of photon-like trigger objects with transverse energy larger than 25
GeV, and the second is to require the smaller threshold of 12 GeV on the photon transverse energy, but at the same time requiring a trigger jet with $E_T$ larger than 20 GeV, and a track with impact parameter larger than 120 microns, reconstructed using the real-time tracker, the SVT. This is a hardware device able to put together in real time discretised information from both the silicon vertex detector (SVX) and the wire chamber (COT), to provide track parameters to the Level 2 trigger. Hits from the COT are connected to form low-resolution tracks, used to get an approximate transverse momentum measurement as well as the $\phi$ position where to look for tracks in the inner SVX. Silicon hit positions as well as the parameters from the low-resolution track are discretised and compared to precalculated patterns stored in an associative memory. Without the need of performing a fit (impossible due to time constraints), this method allows the measurement of track impact parameter with a precision of 47 $\mu$m, out of which 33 coming from the natural beam size; a performance quite similar to the impact parameter resolution of the offline reconstruction.

If we assume that the photon and jet parts of the event are independent (which is the case, since one of the offline requirements is that the jet axis and the photon direction are separated by $\Delta R > 0.7$), the interplay between these two trigger paths can be used to derive the efficiency of the SVT-based trigger, that can therefore directly be measured from data. The analysis is performed first on the unbiased dataset, then on the SVT-based one. Events with photon $E_T$ larger than 25 GeV should be triggered by both, and since the efficiency of the unbiased photon trigger is close to 100% (and anyway quite easy to compute from the MonteCarlo), the percentage of events also selected by the SVT trigger among those with a jet passing offline b-tagging criteria gives a direct measurement of SVT trigger efficiency.

3 Unbiased trigger analysis

Jets are reconstructed using the JetClu clustering algorithm with a cone radius of 0.4, in the rapidity range $|\eta| < 0.7$. The standard CDF b-tagging algorithm SecVtx is applied to search for the presence of a secondary vertex. Events with at least a jet with transverse energy larger than 20 GeV and a positive secondary vertex tag are kept. The secondary vertex tagging algorithm has an efficiency of about 40% for a purity on b jets of about 30%. The tagging efficiency is computer from the MonteCarlo, and cross-checked using a b-enriched sample coming from semileptonic b decays; the purity is directly computed from data fitting the invariant mass of the tracks coming from the reconstructed secondary vertex. If it was possible to reconstruct all the decay products of jet, the mass of b jets would have a distribution peaked at the b mass very well separated from jets coming from charm or light quarks. The presence of neutrals and mis-assigned tracks dilutes quite a lot this separation power, however the vertex mass distribution (figure 1) still has some discrimination power between jets coming from b and lighter quarks. Electromagnetic clusters in the central region ($|\eta| < 1.1$) passing standard photon cuts (isolation, no nearby tracks) and with transverse energy larger than 26 GeV are considered as photon candidates. In the CDF environment there is a huge background coming from high-$E_T$ neutral pions that decay into two photons, with such a small relative angle that they end up in a single electromagnetic cluster. To reduce this background, the CDF electromagnetic calorimeter is equipped with a wire chamber in front of the calorimeter (preshower, CPR) and a second one at the position of the average maximum of the electromagnetic shower (CES). The shape of the signal for each photon candidate is compared to standard shapes taken from a testbeam, and each photon receives a weight correlated to the actual probability of the shower being a photon.

The total cross section is obtained correcting every event with the photon weight from the electromagnetic shower shape, as well as the b purity from the mass of the secondary vertex, and correcting for b-tagging as well as photon identification and trigger efficiency. The jet axis and the photon are required to be isolated by $\Delta R > 0.7$) The resulting cross section as a function of photon $E_T$ is shown in figure 2.
4 Dedicated-trigger analysis

The analysis proceeds as in the previous case, the only difference being that this time since a track with large impact parameter is required at trigger level, trigger efficiency is different from 100% and extracted from data, as the fraction of the events taken by the unbiased trigger and with a jet with positive b-tag, also passing the SVT-based trigger. As seen in figure 3, this trigger efficiency is quite flat as a function of jet E_\text{T}, and around 50%. As seen in figure 4, the result obtained using this dedicated trigger agree well with those of the unbiased analysis for the large photon E_\text{T} region, while extending the photon E_\text{T} range down to 12 GeV.

5 Systematic errors and final results

Two large systematic uncertainties associated to this measurement are common to all jet-based cross section measurements in CDF, ie a 6% uncertainty in the luminosity, and a 4% uncertainty coming to an imprecise knowledge of the jet energy scale. Specific to this analysis are systematics associated to photon and b jet identification. The uncertainty on the photon efficiency is estimated to be 1%, while that associated to photon background to 5%. The comparison of b-tagging efficiency between data and MonteCarlo using b-enriched samples leads to assign a systematic uncertainty of 3% on the efficiency of the algorithm, and the uncertainty on the fraction of b\bar{b} pairs in the jet to another 2%. Extrapolating the trigger efficiency from the high- to the low-photon E_\text{T} region accounts for a 10% uncertainty, mainly of statistical nature due to the small sample from which the extrapolation is performed. But the largest systematic effect comes from the secondary vertex mass fit performed to extract the b purity of the tagged sample. The results are very sensitive to the exact modeling of the
Monte Carlo templates used to fit the data distributions. There are indications that there could be 3% discrepancy between track finding efficiency in data and Monte Carlo. If we translate this value in a ±3% shift of the mass templates, this has an effect on the measured purity of ±20% - 10%. At the moment of writing, no better extraction of this uncertainty is available, but work is going on to provide a more realistic modeling of the effect, namely using measured distributions of track efficiency and account for the difference in tagging requirements between vertices of two or more tracks.

With a total systematic uncertainty of ±24% - 17%, the final measured cross section on 208 pb\(^{-1}\) of CDF data is

\[
\sigma(b + \gamma, |\eta(b)| < 1.5, |\eta(\gamma)| < 1.1, \Delta R(b - \gamma) > 0.7, Et(b) > 20, Et(\gamma) > 12) = \\
90.5 \pm 6.0(stat.)^{+21.7}_{-18.4}(syst.)pb
\]

to be compared to a leading-order prediction from PYTHIA of 69.3 pb. The fact that data are higher than PYTHIA predictions is expected by the absence of higher-order terms in the simulation, and work is going on in obtaining a next-to-leading order prediction for this process.

6 Conclusions

We measured production of photons and b jets using both an unbiased photon trigger and a dedicated trigger that exploits online track information from the SVT to lower the photon Et threshold. The two analyses are in good agreement in the common phase space region, and both show a cross section some 20-30% higher than the leading order QCD predictions, and comparisons with NLO predictions, as well as more detailed systematic studies, are under way.

References

B PRODUCTION AT THE LHC / QCD ASPECTS

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The LHC provides new opportunities to improve our understanding of the $b$ quark using high statistics data samples and the 14 TeV center-of-mass energy. The prospects to measure the cross section for inclusive $b$ production in events containing jets and at least one muon are presented. Studies of detector systematic effects and theoretical uncertainties are included. QCD aspects of the beauty production are discussed.

1 Inclusive $b$-quark production at LHC

1.1. QCD aspects

$B$ production will be one of the most copious sources of hadrons at LHC. Three mechanisms contribute to the beauty production at hadron colliders: gluon-gluon fusion and $q\bar{q}$ annihilation (flavor creation in hard QCD scattering), flavor excitation (semi-hard process) and gluon splitting (soft process). It is important to measure the $B$-hadron $p_T$ spectra within large range to be able to disentangle the contributions of those mechanisms. Flavor creation refers to the lowest-order, two-to-two QCD $b\bar{b}$ production diagrams. Flavor excitation corresponds to diagrams where a $b\bar{b}$ pair from the quark sea of the proton is excited into the final state due to one of the $b$ quarks undergoes a hard QCD interaction with a parton from the other proton. Gluon splitting refers to the processes in which the $b\bar{b}$ pair arises from a $g \rightarrow b\bar{b}$ splitting in the initial or final state. Neither of the quarks from $b\bar{b}$ pair participate in the hard QCD scattering in this case. Inclusive $b$-quark production has been studied at other proton and electron colliders. The observed shapes of
distributions and correlations are reasonably well explained by perturbative QCD. However, the observed cross-sections at the Tevatron (Run I) are larger than QCD predictions\footnote{1,2} which is confirmed by Run II data. Similar effects are observed in $\gamma p$ collisions at HERA\footnote{4,5} and in $\gamma \gamma$ interactions at LEP\footnote{7,8}. The agreement between experiment and theory has improved due to more precise parton density functions and proper estimates of fragmentation effects\footnote{9}. But an agreement is not complete and phenomenological input to the calculations is required.

1.2. $B$ production measurement at LHC

A study\footnote{10} has been performed in CMS on Monte Carlo events generated with PYTHIA\footnote{11} to investigate methods of identifying in CMS of $b$-jets ($b$ “tagging”) in an inclusive sample of events containing jets and at least one muon. We present here the capability to measure the inclusive $b$-quark production cross section as a function of the $B$-hadron transverse momentum $p_T$ and pseudorapidity $\eta$. The study of the CMS capability to measure the inclusive $b$ production is based on full detector simulation. The measurement of the differential cross sections is studied for $B$-hadrons of $p_T > 50$ GeV/c and within the fiducial volume of $|\eta| < 2.4$. The event selection requires a $b$-tagged jet in the fiducial volume to be present in the event. $B$ tagging is based on inclusive secondary vertex reconstruction in jets\footnote{12}. At Level-1 (L1) trigger, the single muon trigger is used. At the High Level Trigger (HLT) we require the “muon + $b$-jet” trigger. The most energetic $B$-hadron inside the phase space defined above is selected. Good correspondence between the generated $B$-particle and the reconstructed $b$-tagged jet is observed. The corresponding relative resolutions for $B$-particles with $p_T > 170$ GeV/c are 13% and 6% for $p_T$ and pseudorapidity, respectively. The average $b$ tagging efficiency is 65% in the barrel region, while the efficiency is about 10% less for the endcap region.

The signal fraction is determined from a fit to the data distribution using the simulated shapes for the signal and background. To do so we apply a lepton tag by selecting inclusive muons. Each reconstructed muon is associated to the most energetic $b$ tagged jet. The muon must be closer to this $b$ tagged jet than to any other jet in the event. Otherwise the event is discarded. In most cases the tagged muon is inside the $b$ jet. The average efficiency of associating the muon with the $b$ tagged jet is 75%. We calculate the transverse momentum of the muon with respect to the $b$-jet axis which effectively discriminates between $b$ events and background. Figure 1 shows an example of the fit of the distribution of the muon $p_T$ with respect to the closest jet, using the expected shapes for the muons from $b$ events, charm events and light quark events. The normalisation of the three contributions are free parameters in the fit. The event fractions are well reproduced within statistical errors\footnote{10}.

![Fig. 1. Fit of the muon $p_T$ spectrum with respect to the closest $b$ tagged jet.](image)

The contributions of tagged muons from $b$ events (dashed curve), $c$ events (dot-dashed curve) and light quark events (dotted curve) as defined by the fit are shown. The solid curve is the sum of the three contributions.
The total event selection efficiency is about 5%. By correcting for the semi-leptonic branching ratio of $b$ quarks and $c$ quarks it amounts to about 25% on average. It turns out that the total efficiency is almost independent of transverse momentum and angle of the $B$-particle. Therefore the measurement of the differential cross section is less affected by systematic uncertainties due to bin-by-bin efficiency corrections.

Several potential sources for systematic uncertainties are considered and their impact on the observed cross section is detailed in \textsuperscript{10}. The largest uncertainty arises from the 3% error \textsuperscript{11} on the jet energy scale which leads to a cross section error of 12% at $E_T > 50$ GeV/c. The estimated statistical, systematic and total uncertainty as function of the $b$ tagged jet transverse momentum with respect to the beam line is shown in Figure 2.

![Figure 2](image.png)

**Fig. 2.** The statistical uncertainty for the cross section measurement (triangles), systematic (squares) uncertainty and total (dots) uncertainty as function of the $b$ tagged jet transverse momentum with respect to the beam line. Total uncertainty comprises the statistical and systematic uncertainties added in quadrature.

The event selection for inclusive $b$ production measurement at CMS will allow to study $b$ production mechanisms on an event sample of 16 million $b$ events for 10 fb\textsuperscript{-1} of integrated luminosity. The $b$ purity of the selected events varies as function of the transverse momentum in a range from 70% to 55%. Our estimate shows that with the CMS detector we can reach 1.5 TeV/c as the highest measured transverse momentum of $B$ hadrons. The results are preliminary, the improvements are likely as further jet calibration tunings, software and analysis algorithm developments are foreseen.

1.3. $\bar{b}b$ correlations

Correlations measurements are foreseen at LHC in order to study details of the production mechanisms discussed in the chapter 1.1. The angular distance $\Delta \phi$ between the $b$ quark directions in the transverse plane is the main discriminating variable to disentangle contributions from the gluon-gluon fusion, gluon splitting and flavor excitation. The $\Delta \phi$ distribution for gluon splitting is slightly peaked at small $\Delta \phi$ values. The angle between the two $b$-quarks produced by the gluon-fusion mechanism has a peak at 180 degrees, as expected, since in the process $gg \rightarrow b\bar{b}$ the $b$ quarks are produced back-to-back in the transverse plane. For the flavor excitation production mechanism the back-to-back topology is preferred too. The MC study with the CMS detector simulation are presented in \textsuperscript{14}.

The ATLAS collaboration is going to use $J/\psi$ from decay of one $B$ particle and semi-leptonic muon decay from another $B$ for the $\bar{b}b$ -correlations measurement \textsuperscript{15}. The $\Delta \phi$ between $J/\psi$ and muon directions distribution is
shown in Figure 3. The muon efficiency as function of $\Delta \Phi_{J/\Psi - \mu}$ is also plotted. The efficiency is rather flat function, which will allow to avoid any significant distortion of the original spectrum.

![Graphs showing muon efficiency as function of $\Delta \Phi_{J/\Psi - \mu}$](image)

Fig. 3. ATLAS study for the $b\bar{b}$-correlations measurement. The left plot shows the $\Delta \phi$ between $J/\Psi$ and muon directions at generator level and after the reconstruction. The right plot shows muon efficiency as function of $\Delta \phi$.

2 Conclusion

All LHC experiments will measure $B$ production cross section at the new accelerator energy frontier, 14 TeV center-of-mass energy. This will provide a new test of the QCD and will fix the normalization for the beauty events background in the new physics searches.

References

Rescattering effects and the determination of the gluon density for $x \ll 1$.

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We consider the possible role of rescattering effects in the determination of the gluon density for the LHC from DIS data. We discuss a method that uses results of $s$-channel calculations to estimate these effects, and comment on potential applications to diffractive and multi-parton interactions.

The Large Hadron Collider will operate with very high gluon luminosities. Production processes initiated by gluons contribute a great many events to a number of cross sections of primary interest for the LHC physics program. Reliable predictions for these cross sections depend on the determination of the gluon density in the proton and its accuracy.\textsuperscript{1,2,3,4,5} As parton luminosities rise steeply for decreasing momentum fraction $x$, a large number of events sample the gluon density at $x \ll 1$.

It has long been known that the theoretical accuracy of gluon-density determinations for $x \ll 1$ is affected by higher-loop $\ln(1/x)$ corrections to QCD evolution equations. See e.g. Ref.\textsuperscript{6} for early numerical investigations. The study of these corrections motivates current searches for evolution schemes (see Refs.\textsuperscript{7,8,9} and references therein) that incorporate the resummation of $\ln(1/x)$ contributions at the next-to-leading-logarithmic accuracy.\textsuperscript{10,11} An improved theoretical control on the $x \ll 1$ region is expected from the inclusion of these terms.

Because the DIS data used at present to extract the gluon density for $x < 10^{-2}$ do not have very high $Q^2$, it is natural to ask whether non-negligible effects on the theoretical accuracy may also come from corrections that are suppressed by powers of $1/Q^2$ but are potentially enhanced as $x \to 0$. These could affect the determination of the gluon density $f_g$ primarily through a contribution $\delta$ to the $Q^2$-derivative of the $F_2$ structure function,

$$\frac{dF_2}{d\ln Q^2} \simeq P_{gq} \otimes f_g \left[ 1 + \delta \right] + \text{quark term} \ , \ \delta \simeq \sum_{k \geq 1} a_k \left( \frac{1}{x} \frac{\Lambda^2}{Q^2} \right)^k .$$ \hspace{1cm} (1)

Here $P_{gq}$ is the perturbative gluon-to-quark evolution kernel, and the correction $\delta$ arises from multi-parton correlation terms in the operator product expansion,

$$F_2 = C \otimes f + \frac{1}{Q^2} C^{(4)} \otimes f^{(4)} + \ldots .$$ \hspace{1cm} (2)

The enhanced $x \to 0$ behavior in Eq. (1) can be produced from graphs with multiple gluon scatterings, and is consistent with observations of approximate geometric scaling in low-$x$ data\textsuperscript{12} and with models for saturation.\textsuperscript{13,14,15}

Standard methods to take account of multiple scatterings are $s$-channel methods (see e.g. the lectures in Ref.\textsuperscript{15}), essentially orthogonal to those of the parton picture. The basic degrees of freedom in the $s$-channel picture are described by correlators of eikonal Wilson lines — at the simplest level, two-point correlators, interpretable as color dipoles. To identify the correction

\textsuperscript{a}Talk presented at the XLII Rencontres de Moriond (La Thuile, March 2007).
from rescattering graphs to the parton result in Eq. (1), a sufficiently precise "dictionary" is needed to connect the two pictures. Ref.\textsuperscript{16} presents an approach to analyze this connection.

The method is based on constructing explicitly an $s$-channel representation for the renormalized parton distribution function in terms of Wilson-line matrix elements, convoluted with lightcone wave functions. In this representation the quark distribution $f_q$ is given by the coordinate-space convolution

$$x f_q(x, \mu) = \int dx \int db \ u(\mu, z) \ \Xi(z, b) ,$$  \hspace{1cm} (3)

where $\Xi$ is the hadronic matrix element of eikonal-line operators,

$$\Xi(z, b) = \int [dP] \langle P| P' \frac{1}{N_c} \text{Tr}\{1 - V^\dagger(b + z/2) V(b - z/2)\}|P\rangle ,$$  \hspace{1cm} (4)

$$V(z) = \mathcal{P} \exp \left\{ -ig \int_{-\infty}^{+\infty} dz^- A^+_{\mu}(0, z^-, z) t_\mu \right\} ,$$

$z$ is the transverse separation between the eikonal lines, $b$ is the impact parameter, and the function $u(\mu, z)$ is evaluated explicitly in Ref.\textsuperscript{16} at one loop using the $\overline{\text{MS}}$ scheme for the renormalization of the ultraviolet divergences $z \to 0$.

The representation (3), once evaluated in a well-prescribed renormalization scheme, is the key ingredient that allows one to relate\textsuperscript{16,17} results of $s$-channel calculations for structure functions to the OPE factorization (2), and, in particular, to identify the power-suppressed corrections arising from the multiple gluon scatterings, treated in the high-energy approximation of Eq. (4). These corrections are found to depend on moments of $\Xi$, schematically in the form

$$\frac{dF_2}{dlnQ^2} = \left( \frac{dF_2}{dlnQ^2} \right)_{LP} + \sum_{n=1}^{\infty} R_n \frac{\lambda^2(n)}{(Q^2)^n} ,$$  \hspace{1cm} (5)

where the first term in the right hand side is the leading-power parton result, and the $\lambda^2$ in the subleading terms are given by the analytically continued moments

$$\lambda^2(-v) = \frac{1}{\Gamma(v)} \int \frac{dz}{\pi z^2} (z^2)^{v-1} \int db \ \Xi(z, b) .$$  \hspace{1cm} (6)

The coefficients $R_n$ are evaluated to order $\as^3$, as functions of $x$ and $\ln Q^2$, from the lightcone wave functions, while the moments $\lambda^2(n)$ are dimensionful nonperturbative parameters, to be determined from comparison with experimental data.

In practice, the usefulness of the result in Eqs. (5),(6) comes from the fact that the hadronic matrix element $\Xi$ can be related by a short-distance expansion for $z \to 0$ to a well-prescribed integral of the gluon distribution function\textsuperscript{16}. Then the moments $\lambda^2$ can be parameterized in terms of the factorization/renormalization scales at which the gluon distribution and the strong coupling are evaluated. These scales are to be taken of the order of the inverse mean transverse distance $1/|z|$, and can be tuned to the data. The result of doing this for $F_2$ data\textsuperscript{18} at both low and high $Q^2$ is shown in Fig. 1 in the left hand side plot.\textsuperscript{17} The corresponding power correction in Eq. (5) is plotted on the right hand side of Fig. 1. Here we show the correction normalized to the full answer and multiplied by $-1$, using CTEQ parton distributions.\textsuperscript{19}

Fig. 1 indicates that with physically natural choices of the parameters in the nonperturbative matrix elements (4) one can achieve a sensible description of data for $x < 10^{-2}$ in a wide range of $Q^2$ and still have moderate power corrections to $dF_2/d\ln Q^2$. Corrections turn out to be negative and below 20 % for $x \gtrsim 10^{-4}$ and $Q^2 \gtrsim 1$ GeV$^2$. However, Fig. 1 also indicates that for small $x$ the corrections fall off slowly in the region of medium $Q^2$, $Q^2 \simeq 1 - 10$ GeV$^2$, behaving effectively like $1/Q^4$ with $\lambda$ close to 1.\textsuperscript{17} For instance, one has $\lambda \simeq 1.2$ for the curve $x = 10^{-3}$...
in the right-hand side plot of Fig. 1. As a consequence of the slowly decreasing behavior, the power corrections stay on the order of 10% up to $Q^2$ of a few GeV$^2$ for $x \lesssim 10^{-3}$. This slow fall-off differs from parameterizations of higher twist commonly used in global analyses (see e.g. Ref. 3), and may be relevant for phenomenology as it affects the medium $Q^2$ region of the data used to extract $f_g$ at low $x$.

Corrections larger than for $F_2$ are found for the longitudinal component $F_L$. This provides additional motivation for the forthcoming $F_L$ measurements$^{20}$, as well as fits$^{21}$ investigating power-like terms in $F_L$ (at both high and low $x$). We observe also that Fig. 1 is obtained using NLO parton distributions, and the decrease in the low-$x$ gluon at NNLO$^2$ could be consistent with the possibility that NNLO parton distributions correspond to smaller power corrections. However, the detailed interpretation of this behavior will be subtle, as distinctly different dynamics drive the power-like and NNLO effects, unlike the high-$x$ case in the analyses$^{3,22}$.

It is worth emphasizing that the above results depend on the validity of the high-energy approximation and s-channel representation (3), and the perturbation expansion for $u(\mu, z)$. The rationale for this expansion lies with the dynamical cut-off on large distances $z$ imposed by unitarity requirements ("black disc" limit) on the correlator $\Xi$. But the size of this cut-off at collider energies is difficult to determine. The highest sensitivity to it may come from measurements of the diffractive part of the DIS cross section. In this case the s-channel representation is bilinear in $\Xi$. The comparison of diffractive data with theoretical predictions based on diffractive parton distributions indicates that the dynamical cut-off lies at substantially higher momenta for color-octet eikonal-line matrix elements than for color-triplet. That is, gluons' shadow is stronger. See e.g. Ref. 27 for a recent discussion. In diffractive DIS this can be linked$^{23,28}$ to the distinctive pattern of the observed scaling violation$^{24}$ and detailed features of the associated jet distributions$^{25,29,30}$. More generally, it suggests that the expansion used is better justified for processes directly coupled to the gluon distribution than for $F_2$, see e.g. applications to $F_L$ (or its diffractive component$^{31}$) and jet final states. A critical discussion, including the quark case, is given in Ref. 16.

Note that the question of how to perform QCD calculations that incorporate multiple scatterings along with perturbative evolution becomes especially compelling in the case of Monte Carlo event generators$^{32}$. The application considered above deals with corrections to $Q^2$ evolution, i.e., a picture based on strongly ordered $k_\perp$'s. But the method relies on high-energy approximations that may be better suited for extension to evolution with ordering in energies, or angles. It could thus more likely be adapted to the modeling of multiparton processes in Monte Carlo generators$^{33,34}$ based on high-energy evolution equations.
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32. See e.g. talks by S. Gieske, T. Sjöstrand, R. Engel at Workshop on Multiparton Interactions, DESY, May 2007.
HADRONIC FINAL STATES AND SPECTROSCOPY
IN EP COLLISIONS AT HERA

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Recent results on spectroscopy and the measurement of hadronic final states in ep collisions from the H1 and ZEUS collaborations are presented using data sets with an integrated luminosity between 44 and 121 pb$^{-1}$ collected during the HERA I running period. Besides a search for resonant states which could be interpreted as pentaquarks, a study of charged particle momentum spectra in the Breit frame and a measurement of neutral strange hadron production is shown. Furthermore two recent measurements of prompt photons are presented and compared with pQCD calculations. The measurements are performed in photoproduction ($\gamma p$) with a four-momentum transfer squared $Q^2 \sim 0$ GeV$^2$ or in deep inelastic scattering (DIS) at $Q^2 > 1$ GeV$^2$.

**Keywords:** pentaquark, prompt photon, charged particle multiplicity, spectroscopy, HERA.

1 Spectroscopy

1.1 Measurement of $K^0_S$, $\Lambda$ and $\bar{\Lambda}$ Production

The production of the neutral strange hadrons $K^0_S$, $\Lambda$ and $\bar{\Lambda}$ has been measured$^1$ by the ZEUS collaboration. In addition to differential cross sections, further measurements are presented, such as the baryon-antibaryon asymmetry, baryon-to-meson ratio, ratio of strange-to-light hadrons, and the $\Lambda$ ($\bar{\Lambda}$) transverse spin polarisation.

The ARIADNE Monte Carlo program can describe the differential cross sections in transverse momentum $P_T$ and pseudorapidity $\eta = -\ln \tan(\theta/2)$ of the hadrons reasonably well when adjusting the strangeness-suppression factor to $\lambda_s = 0.3$, although the cross section at high $Q^2$ is overestimated. ARIADNE adjusted to $\lambda_s = 0.22$ is giving less satisfactory results.
The baryon-to-meson ratio $R = (N(\Lambda) + N(\bar{\Lambda}))/N(K_S^0)$ varies between 0.2 and 0.5 in DIS, which is in agreement with measurements at $e^+e^-$ colliders, where $R$ ranges between 0.2 and 0.4. $R$ is described by ARIADNE ($\lambda_s = 0.3$) to better than 10–20%. In $\gamma p$ a dijet sample is compared to the PYTHIA event generator. At large values of $x^\gamma_{\text{HIS}} > 0.75$, where $x^\gamma_{\text{HIS}}$ is a measure of the fraction of the photon energy transferred to the dijet system, $R$ is found to be roughly 0.4, which is in agreement with the PYTHIA prediction and also corresponds to the values seen in DIS as discussed above. For low values of $x^\gamma_{\text{HIS}}$, $R$ rises to a value of about 0.7, while PYTHIA predicts a flat $x^\gamma_{\text{HIS}}$ dependence.

The measured ratio of strange-to-light hadrons $T = N(K_S^0)/[N(\pi^+)/N(K^+) + N(\bar{p}) + N(\bar{\Lambda})]$ lies between 0.05 and 0.1 varying with $p_T$. A comparison with ARIADNE suggests a strangeness-suppression factor $\lambda_s < 0.3$.

1.2 Search for Baryonic States Decaying to $\Xi\pi$ in Deep Inelastic Scattering

A search for narrow baryonic resonances decaying into $\Xi^{-}\pi^{-}$ or $\Xi^{-}\pi^{+}$ and their antiparticles is carried out$^3$ with the H1 detector in DIS. No signal is observed for a new baryonic state in the mass range 1600–2300 MeV in neither the doubly charged ($\Xi^{-}\pi^{-}, \Xi^{+}\pi^{+}$) nor the neutral ($\Xi^{-}\pi^{0}, \Xi^{+}\pi^{-}$) decay channels. In the neutral charged combinations there is a clear signal of the well-known $\Xi(1530)^0$ resonance with 170 signal events. The NA49 collaboration$^4$ observed two baryonic resonances with masses around 1.86 GeV, which can be interpreted as pentaquark states. The observation cannot be confirmed by the H1 measurement.

Mass dependent upper limits at 95% confidence level relative to the $\Xi(1530)^0$ signal $R_{\Xi,1}(M)$ are derived making use of a modified frequentist approach. The limits vary between 0.15 and 0.6 with values of $R_{\Xi,1}(1860) \sim 0.2(0.5)$ for the doubly (neutral) charged combinations. The sum of all charged combinations yields an upper limit of $R_{\Xi,1}(1860) \sim 0.5$, which is in agreement with the upper limit of 0.29 derived by the ZEUS collaboration$^5$.

1.3 Charged Particle Production in high $Q^2$ Deep Inelastic Scattering

The process of parton fragmentation and hadronisation has been studied$^6$ by H1 using charged particle momentum spectra at high $Q^2$. The measurement is performed in the current region of the Breit frame with an energy scale given by $Q/2$. Observables in the current region of the Breit frame can be compared with similar observables measured in one hemisphere of $e^+e^-$ annihilation events, where the energy scale is half the centre-of-mass energy $E^*/2$. In the Breit frame the scaled momentum variable $x_p$ is defined as $p^+_{h}/(Q/2)$, where $p^+_{h}$ is the momentum of a charged track. In $e^+e^-$ annihilations the corresponding observable is $2p^+_{h}/E^*$.

The inclusive, event normalised charged particle scaled momentum spectrum defined as $D(x_p,Q) = (1/N_{\text{event}})dn/dx_p$ is shown in Fig. 1 as a function of $Q$ for different bins of $x_p$ and is compared with results from $e^+e^-$ annihilation events. The $ep$ and $e^+e^-$ data are in broad agreement supporting the concept of quark fragmentation universality. The spectra become softer when moving from low to high $Q$, which implies a substantial increase in the number of hadrons with a small share of the initial parton’s momentum. These scaling violations are assumed to be caused by parton splitting in QCD, i.e. the same effect that causes the scaling violations observed in deep inelastic structure functions. The parton shower model as implemented in RAPGAP gives the best description of the charged particle momentum spectra over the full range of $x_p$.

1.4 Measurement of $\bar{d}$ and $\bar{p}$ Production in Deep Inelastic Scattering

The production of (anti)deuterons, $d(\bar{d})$, and (anti)protons, $p(\bar{p})$, has been studied$^7$ in DIS with the ZEUS detector and represents the first measurement of $d$ production in DIS. The
(anti)deuterons and (anti)protons are identified by means of the energy-loss measurement $dE/dx$ in the central tracking detector.

The corrected $d$ production rate is found to be 3–4 orders of magnitude lower than the corrected $p$ yield, which is in agreement with the H1 published data\(^8\) in $\gamma p$. Furthermore the measured $p/p$ ratio is consistent with unity. Within the given statistics antitritons have not been observed.

2 Prompt Photon Production

2.1 Inclusive Prompt Photon Production in Deep Inelastic Scattering

A measurement of prompt photons in DIS\(^9\) has been presented by the H1 collaboration. Compared to the previous measurement\(^10\) of prompt photons in DIS, the total cross section expectation is increased by roughly a factor of 10 due to a markedly extended phase space. The photon's transverse energy and pseudorapidity range is given by $3 < E_T^\gamma < 10$ GeV and $-1.2 < \eta^\gamma < 1.8$.

Photons are identified by a compact electromagnetic cluster in the main calorimeter with no track pointing to it. The major background due to neutral hadrons inducing multi-photon clusters is considerably reduced by an infrared-safe isolation criteria. The extraction of the photon content from the remaining neutral hadron background is done by a fit to the output of a multivariate shower shape analysis. The measured differential cross sections $d\sigma/dE_T^\gamma$ and $d\sigma/d\eta^\gamma$ are compared to a LO($\alpha^3$) calculation\(^11\). The differential cross section $d\sigma/dE_T^\gamma$ is well described in shape, though the calculation is too low in normalisation. The underestimation is most visible at central pseudorapidities as can be seen in Fig. 2 a).

At large pseudorapidities the dominant contribution is photon radiation by the quark, while at low pseudorapidities, close to the scattered electron, the contribution of photon radiation from the electron is dominant.

2.2 Measurement of Prompt Photons with Associated Jets in Photoproduction

The photoproduction of prompt photons, together with a separate jet in addition to the photon, has been studied\(^12\) with the ZEUS detector at HERA.

The photon identification is based on the conversion probability of photons to $e^+e^-$ in front
Figure 2: Inclusive prompt photon differential cross section $d\sigma/d\eta^\gamma$ in DIS as measured by H1 (a) and the photon plus accompanying jet differential cross section $d\sigma/dE_T^\gamma$ in $\gamma p$ as measured by the ZEUS collaboration (b). The measurements are compared to various pQCD calculations.

of a preshower detector (RPQ). The photon kinematics is restricted to $5 < E_T^\gamma < 16$ GeV and $-0.7 < \eta^\gamma < 1.1$, while the accompanying jet is selected in the kinematic range $6 < E_T^{jet} < 17$ GeV and $-1.6 < \eta^{jet} < 2.4$. A similar infrared-safe isolation criteria as in the above H1 analysis is used. Differential cross sections as functions of $E_T^\gamma$ and $\eta^\gamma$ are compared to two NLO calculations (cf. references in$^{12}$). As visible in Fig. 2 b) both calculations describe the data rather well, however underestimate the data at low $E_T^\gamma$. Another calculation (cf. reference in$^{12}$) based on a $k_T$ factorisation approach yields the best description of the cross sections, particularly at low $E_T^\gamma$. When raising the minimum transverse energy to $E_T^\gamma > 7$ GeV, the data is well described by all three predictions.

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A GENERAL SEARCH FOR NEW PHENOMENA IN $e^-p$ SCATTERING AT HERA

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A model-independent search for deviations from the Standard Model prediction is performed in $e^-p$ collisions at HERA II using H1 data recorded during the years 2005 and 2006, corresponding to an integrated luminosity of 159 pb$^{-1}$. All event topologies involving isolated electrons, photons, muons, neutrinos and jets with high transverse momenta are investigated in a single analysis. Events are assigned to exclusive classes according to their final state. A statistical algorithm is used to search for deviations from the Standard Model in the distributions of the scalar sum of transverse momenta or invariant mass of final state particles and to quantify their significance. A good agreement with the Standard Model prediction is observed in most of the event classes. No significant deviation is found in the phase-space and event topologies covered by this analysis.

1 Introduction

At HERA electrons$^a$ and protons collide at a centre-of-mass energy of up to 319 GeV. These high-energy electron-proton interactions provide a testing ground for the Standard Model (SM) complementary to $e^+e^-$ and $pp$ scattering.

The approach described in this paper$^1$ closely follows the strategy of the previously published H1 analysis using HERA I data$^2$. It consists of a comprehensive and generic search for deviations from the SM prediction at large transverse momenta. The analysis covers phase-space regions where the SM prediction is sufficiently precise to detect anomalies and does not rely on assumptions concerning the characteristics of any SM extension. Using the complete HERA II $e^-p$ data sample, this is the first general search performed on a large data set from electron-proton collisions.

$^a$In this paper “electrons” refers to both electrons and positrons, if not otherwise stated.
Figure 1: The data and the SM expectation for all event classes with observed data events or a SM expectation greater than 1 event. The analysed data sample corresponds to an integrated luminosity of 159 pb$^{-1}$. The error bands on the predictions include model uncertainties and experimental systematic errors added in quadrature.

2 Data analysis

The event sample studied consists of the full 2005–2006 HERA II $e^-p$ data set, corresponding to an integrated luminosity of 159 pb$^{-1}$. All final states with at least two objects with $P_T > 20$ GeV in the polar angle range $10^\circ < \theta < 140^\circ$ are investigated. Considered objects are electrons ($e$), photons ($\gamma$), muons ($\mu$), jets ($j$) and neutrinos ($\nu$) (or non-interacting particles). The identification criteria for each type of object are similar to those applied in the published HERA I analysis, ensuring an unambiguous identification while keeping high efficiencies. All objects are required to be isolated from each other by a minimum distance $R$ of 1 unit in the $\eta - \phi$ plane. The events are classified into exclusive event classes according to the number and types of objects. This exclusive classification ensures a clear separation of the final states and allows an unambiguous statistical interpretation.

As this analysis investigates all final state topologies of $ep$ interactions at high transverse momentum, a precise and reliable estimate of all relevant HERA processes is needed. Hence, several Monte Carlo generators are used to generate a large number of events in all event classes, carefully avoiding double-counting of processes. The simulation contains the order $\alpha_S$ matrix elements for QCD processes, while second order $\alpha$ matrix elements are used to calculate QED processes. Additional jets are modeled using leading logarithmic parton showers as representation of higher order QCD radiation. All processes are generated with a luminosity significantly higher than that of the data.

The results of the data analysis are summarised in figure 1, which presents the event yields subdivided into event classes for the data and SM expectation. All event classes with observed
Figure 2: The $-\log \hat{P}$ values for the data event classes and the expected distribution from MC experiments as derived by investigating the $\sum P_T$ distributions (left) and $M_{\text{all}}$ distributions (right) with the search algorithm.

data events or a SM prediction greater than 1 event are shown\(^5\). In each class, a good agreement between the number of observed data events and the SM prediction is seen.

No data events are observed in the event classes $\mu-j-\nu$ and $e-j-\nu$ where the largest discrepancy between the data and the SM prediction was found in the analysis of the HERA I data (see section 4). These classes correspond mainly to high $P_T$ $W$ production with subsequent leptonic decay, where deviations in the $e^+p$ data continue to be observed\(^3\). The total SM expectation amounts to $1.2 \pm 0.2$ and $2.5 \pm 0.8$ in the $\mu-j-\nu$ and $e-j-\nu$ classes, respectively.

3 Search for deviations

In order to quantify the level of agreement between the data and SM expectation and to identify regions of possible deviations, the invariant mass $M_{\text{all}}$ and sum of transverse momenta $\sum P_T$ distributions of all reliable event classes are systematically investigated using the same search algorithm as developed for the previous publication. A region is defined as a sample of connected histogram bins, which have at least the size of twice the resolution of the observable. A statistical estimator $\hat{P}$ is defined to determine the region of most interest by calculating the probability that the SM expectation fluctuates upwards or downwards to the data. This estimator is derived from the convolution of a Poisson probability density function (pdf) to account for statistical errors with a Gaussian pdf to include systematic uncertainties. A possible sign of new physics is found if the expectation significantly disagrees with the data, and thus the region of most interest (greatest deviation) is given by the region having the smallest $p$-value, $p_{\text{min}}$. This method finds narrow resonances and single outstanding events, as well as signals spread over large regions of phase-space in distributions of any shape.

The fact that somewhere in the distribution a fluctuation with a value $p_{\text{min}}$ might occur is taken into account by calculating the probability $\hat{P}$ to observe a deviation with a $p$-value $p_{\text{min}}$ at any position in the distribution. Thus $\hat{P}$ is the central measure of significance of the deviation found. To determine $\hat{P}$, hypothetical data histograms are diced according to the pdf of the expectation. The value of $\hat{P}$ is then defined as the fraction of hypothetical data histograms with a $p_{\text{min}}$-value smaller than the $p_{\text{min}}$-value obtained from the data, and consequently the event class of most interest for a search is the one with the smallest $\hat{P}$-value.

The overall level of agreement between the data and SM expectation can be quantified further by taking into account the large number of event classes studied in this analysis. Among all

\(^5\)The $\mu-\nu$ event class is discarded from the present analysis. It is dominated by events in which a poorly reconstructed muon gives rise to missing transverse momentum, taking the neutrino signature.
classes, there is some chance that small \( \hat{P} \) values occur. This probability can be calculated by replacing all data distributions by hypothetical Monte Carlo (MC) distributions based on the SM expectation. The complete statistical algorithm is applied on this MC experiment, representing a single HERA experiment with an integrated luminosity of 159 pb\(^{-1}\). The expectation for the \( \hat{P} \) values of the data is then given by the distribution of \( \hat{P}_{SM} \) values derived from many MC experiments. In the case that deviations arise from statistical or systematical fluctuations only, the distribution of \( \hat{P} \)-values obtained from data and MC experiments are compatible.

The results of the search for deviations between data and SM expectation are summarised in figure 2. Shown is the negative base-10 logarithm of the \( \hat{P} \) values obtained from the real data compared to the expectation derived from a large set of MC experiments. The comparison is presented separately for the scans of the \( M_{\text{all}} \) and \( \sum P_T \) distributions. All \( \hat{P} \) values range from 0.01 to 0.99, corresponding to event classes where no significant discrepancy between data and SM expectation is observed. These results are in good agreement with the expectation from MC experiments.

Although data events are observed in the \( j-j-j-j \) and \( j-j-j-j-\nu \) event classes, no reliable \( \hat{P} \) values can be calculated for these classes due to uncertainties of the SM prediction at highest \( M_{\text{all}} \) and \( \sum P_T \) values\(^2\). These event classes are not considered to search for deviations from the SM in this extreme kinematic domain. Consequently, these event classes are not taken into account to determine the overall degree of agreement between the data and the SM.

4 Comparison with HERA I results

While good agreement in all event classes is observed between the data and the SM prediction in the present HERA II analysis, some discrepancy is found in the previously published general search on the HERA I data. Complementary to the pure \( e^- p \) event sample studied here, the HERA I data set is largely dominated by positron-proton collisions. There the most significant deviation is found in the \( \mu-j-\nu \) event class with \( \hat{P} \) values of 0.01 and 0.001 for the scan of the \( M_{\text{all}} \) and \( \sum P_T \) distributions, respectively. The global probability to find at least one event class with a \( \hat{P} \) value smaller than that of the \( \mu-j-\nu \) class in the HERA I data amounts to 28% for the \( M_{\text{all}} \) and 3% for the \( \sum P_T \) distributions.

5 Conclusions

The data collected with the H1 experiment during the years 2005–2006 (HERA II) have been investigated for deviations from the SM prediction at high transverse momentum. All event topologies involving isolated electrons, photons, muons, neutrinos and jets are investigated in a single analysis. This is the first general search performed on a large set of data from electron–proton collisions. A good agreement between data and SM expectation is found in most event classes. In each event class the invariant mass and sum of transverse momenta distributions of particles have been systematically searched for deviations using a statistical algorithm. No significant deviation from the SM is observed in the phase–space and event topologies covered by this analysis.

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New trends in modern event generators

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Abstract

Some features of modern simulation tools for high-energy physics are reviewed.

1 Introduction: The next generation of event generators

In the past decades, event generators have become increasingly important for the planning of collider experiments and analyses of their data. In the LHC era this trend will become even more pronounced, since many of the interesting signals expected at the LHC - such as signals for the Higgs boson or alternative electroweak symmetry breaking mechanisms, supersymmetry, etc. - are severely hampered by large backgrounds, with a significant influence of QCD. Thus, the success of the LHC probably rests on a precise understanding of these backgrounds. Examples for this include the effect of the central jet veto in vector boson fusion, producing the Higgs boson, and in multi-jet backgrounds to SUSY searches. In view of this, it is obvious that many of the old tools need to be replaced by newer, and better ones, such as Pythia8 [1], Herwig++ [2], and SHERPA [3]. Their ongoing construction in fact reflects increased experimental needs. In many cases, they therefore incorporate new, better simulation methods, many of which are connected to the systematic inclusion of higher-order QCD-corrections.

2 Parton level: Calculation of signal and background

Many of the apparent improvements of current event generators rely on the inclusion of higher-order corrections; one method is to use multi-leg tree-level matrix elements (MEs) as a base for simulation. There are a number of such tools on the market, which either evaluate Feynman graphs using the helicity method (for instance [4]) or recursion relations, e.g. [5].
Table 1: Performance of different calculational methods for multi-leg QCD matrix elements. Time is given for the calculation of $310,000$ phase space points. Clearly, the CSW recursion relations (here labeled with MHV) are superior in performance; the fact that only up to two quark lines have been included in the respective algorithm, has negligible influence on the final result.

<table>
<thead>
<tr>
<th>Process</th>
<th>Cross section</th>
<th>time (helicity)</th>
<th>time (MHV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$g g \rightarrow j j$</td>
<td>$745.85 \mu b \pm 0.10%$</td>
<td>$66 s$</td>
<td>$44 s$</td>
</tr>
<tr>
<td>$g g \rightarrow j j j$</td>
<td>$81.274 \mu b \pm 0.20%$</td>
<td>$1400 s$</td>
<td>$166 s$</td>
</tr>
<tr>
<td>$g g \rightarrow g g g g$</td>
<td>$10.145 \mu b \pm 0.23%$</td>
<td>$90 ks$</td>
<td>$0.6 ks$</td>
</tr>
<tr>
<td>$j j \rightarrow j j j j j$</td>
<td>$23.208 \mu b \pm 0.26%$</td>
<td>$210 ks$</td>
<td>$5.8 ks$</td>
</tr>
<tr>
<td>$g g \rightarrow g g g j j$</td>
<td>$2.6915 \mu b \pm 0.15%$</td>
<td>$- $</td>
<td>$17 ks$</td>
</tr>
<tr>
<td>$j j \rightarrow j j j j j j$</td>
<td>$7.3294 \mu b \pm 0.17%$</td>
<td>$- $</td>
<td>$122 ks$</td>
</tr>
</tbody>
</table>

Table 2: Unweighting efficiencies obtained with the integration in AMEGIC++. 

<table>
<thead>
<tr>
<th>$g g \rightarrow g g$</th>
<th>$j j \rightarrow j j$</th>
<th>$g g \rightarrow g g g g$</th>
<th>$j j \rightarrow j j j j j j$</th>
<th>$g g \rightarrow g g g j j$</th>
<th>$j j \rightarrow j j j j j j j j$</th>
</tr>
</thead>
<tbody>
<tr>
<td>5.8%</td>
<td>1.6%</td>
<td>2.0%</td>
<td>0.5%</td>
<td>0.9%</td>
<td>0.2%</td>
</tr>
</tbody>
</table>

However, apart from being able to calculate the MEs quickly, also the ability to integrate efficiently over the final state particles' phase space is a major obstacle for a satisfying performance of such tools. In the following, some results addressing both issues, are presented, serving as illustrative examples for current performances. In Tab. 1, the performance of the helicity method is compared to the performance of the CSW recursion relations [6]. For pure QCD processes both approaches have been implemented in the parton level generator AMEGIC++ [7], which is a central part of SHERPA. The phase space integration and the corresponding unweighting efficiencies rely on an integrator based on a hierarchical antenna generation (HAAG) [8], which has been further improved with VEGAS [9]. Results for the unweighting efficiencies are displayed in Tab. 2.

3 From parton to hadron level

For experimental analyses, however, parton level results, discussed in the previous section, are of limited interest only. This is due to the fact that the experimental discussion of jets is based on hadrons rather than on partons. At the moment, the transition from partons to hadrons can be described with phenomenological models only, which depend on tunable parameters. In order to guarantee the validity of such a tuned parameter set, the partons entering these models should have comparable distances in phase space. Ultimately, this is what the parton shower (PS), modelling secondary Bremsstrahlung emissions, is responsible for. When comparing matrix elements (MEs) with the PS, it becomes apparent that they perform best in different regimes of particle creation. While MEs essentially are well-suited to describe hard, large-angle emissions, taking interferences into account, the PS covers especially soft and collinear emissions, resumming corresponding large logarithms. It is therefore natural to try to combine both descriptions into a unified one, employing the best of both approaches for an improved simulation. The catch in so doing is to avoid double-counting of emissions into the same region of phase space and to preserve the correct treatment of leading logarithms. An algorithm satisfying these requirements has been presented in [10] for the case of $e^+ e^- \rightarrow$ hadrons, where its accuracy up to next-to leading logarithmic order has been proven. An extension to hadronic collisions has been presented in [11]. This algorithm aims at a description of all jet emissions correct at tree level plus leading logarithms, with all soft and collinear emissions correctly taken care of in the PS. To achieve this, parton emission is separated into two regimes, one for jet production and one for jet evolution, through a $k_t$-algorithm [12]. Jets are then produced according to tree level MEs, the corresponding configurations are re-weighted with analytical Sudakov form factors and running $\alpha_s$ weights. In the PS, production of additional hard jets is vetoed. Altogether, this algorithm has been implemented
in a process-independent way in SHERPA [3], allowing for careful validation and cross checks with experimental data or other calculations [13].

As a non-trivial check of the quality in describing the QCD radiation pattern through the merging approach, consider the azimuthal decorrelation of jets in $pp$ collisions at the Tevatron, Run II, as presented in [14]. This observable effectively tests additional radiation, both hard and soft, in inclusive QCD dijet production. The agreement of the results of a SHERPA simulation with the experimental results is remarkable, cf. the upper left panel of Fig. 1. A prime example for the predictive power of the merging approach of SHERPA is the case of inclusive $Z$ as measured by the D0 collaboration at Tevatron, Run II [13]. There, SHERPA is not only perfectly capable to predict the relative multiplicities of associated jets, cf. the upper right panel of Fig. 1; it also yields an improved description of the jet kinematics. This is illustrated in the lower panels of Fig. 1, where the transverse momentum distribution of the third-hardest jet (left) and the azimuthal correlation of the two leading jets (right) are displayed. In so doing, its abilities stretch beyond those of other, more traditional event generators, which do not rely on such a merging approach.

4 Modelling hadron decays

Another improvement of modern event generators when compared to traditional ones rests in the description of hadron decays and decay chains. Apart from the inclusion of spin correlations [16], modelling the effect of interferences in decay chains, apparent refinement of the simulation can be achieved by using better form factor models in decays, leading to non-flat phase space distributions and by an upgraded description of mixing effects, like, e.g. $BB$ mixing. Some of these refinements are exemplified in Fig. 2. There, in the left panel, the effect of different form factor models [17] on $m_{\pi\pi}$ in decays $\tau \rightarrow \pi\pi\nu_{\tau}$ are compared with experimental data from [18], whereas in the right panel the asymmetry of $J/\psi K_{S}$ final states in $B$ decays is displayed.
5 Conclusions

In this contribution, the need for new simulation tools in preparation for a successful LHC era has been motivated. These new tools become mandatory due to the abundance of backgrounds, shadowing potentially interesting signals. An apparent feature of modern event generators, improving traditional ones, rests in the systematic inclusion of higher-order QCD corrections through merging or matching algorithms. One of them has been shortly discussed, and results obtained with it have been presented. In order to realise tree-level merging algorithms, multi-log tree level parton level event generators are an important ingredient, and some recent developments concerning the efficient calculation of corresponding cross sections have been shown. Finally, another rectification included in modern tools, consists in a better understanding and modelling of hadron (especially B and D) and τ decays and in the simulation of non-trivial quantum interference effects.

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Herwig++

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We describe the recent development of the Herwig++ event generator.

Keywords: Hadron Colliders, Monte Carlo Simulations

1 Introduction

Monte Carlo event generators have become an essential part of all experimental analyses in particle physics. While it has been possible to extend and improve the existing HERWIG\textsuperscript{1} program over many years making further improvements is increasingly difficult. In order to include our improved understanding of the physics and recent theoretical developments a programme is therefore underway to produce a new simulation in C++, Herwig++\textsuperscript{2,3}, based on the same physics philosophy and models. This is part of a wider program within the Monte Carlo community to produce a new generation of simulations for the LHC.\textsuperscript{4}

We will present the recent physics developments in the Herwig++ simulation, concentrating on improvements to the simulation of QCD radiation and new physics, and plans for further improvements.

2 Simulation of QCD Radiation

The main change between the HERWIG and Herwig++ programs is in the simulation of perturbative QCD radiation. While both programs use an angular-ordered parton shower designed to treat soft-gluon interference effects the algorithm in Herwig++\textsuperscript{2,5} has a number of improvements: invariance under boosts along the jet axis; improved treatment of radiation from heavy quarks; and better coverage of the soft region of phase space.
In particular the new algorithm uses an improved evolution variable and the quasi-collinear splitting functions\(^6\) to give better treatment of the radiation from massive particles. In FORTRAN HERWIG the "dead-cone" approximation of forbidding radiation with angle less than \(m/E\), where \(m\) is the mass and \(E\) the energy of the heavy particle, was used. In Herwig++ this is replaced by a smooth suppression of radiation in the direction of the particle.

The improved treatment of the kinematics of the branchings in the shower means that the soft-region of phase space in \(e^+e^- \rightarrow q\bar{q}\) is smoothly covered with radiation from the quark and anti-quark filling separate regions of phase space which cover the whole region for soft emission without overlapping, as was the case with the FORTRAN algorithm.

A major new feature is the inclusion of radiation from the decaying particle in heavy particle decays, for example \(t \rightarrow bW^+\). This means that in these decays the soft region for gluon emission is fully covered, whereas in the FORTRAN program, which only included radiation from the decay products, part of the soft region was not filled.\(^7\) This makes correcting the parton shower using the exact single emission matrix element simpler. These corrections are now included for \(e^+e^- \rightarrow q\bar{q}\), top decay\(^8\) and the Drell-Yan process.

Another key feature of the new algorithm is that momentum reshuffling needed to ensure energy and momentum conservation is under greater analytic control which will make it easier to match with higher order matrix elements. A number of developments in this area are underway.\(^9\)

3 BSM Physics

The existing HERWIG program includes a detailed simulation of supersymmetric (SUSY) models\(^10\) including both spin correlation effects\(^11\) and R-parity violating models.\(^12\) However, while the simulation of SUSY models was highly sophisticated, extending the simulation to other models of new physics was difficult and time consuming. In the new simulation we have adopted an entirely different approach for the inclusion of new physics models.\(^13\) In the FORTRAN simulation the matrix element for each new scattering process and decay was added by hand. In the new simulation we have included a library based on the HELAS\(^14\) formalism which is used in all matrix element calculations. The spin structures for the possible \(2 \rightarrow 2\) matrix elements and \(1 \rightarrow 2\) decays are included, based on the possible Feynman diagrams for each combination of the spins of external particles. The possible scattering processes and decays are then
automatically calculated from the Feynman rules implemented in the code. Using the HELAS formalism allows us to include spin correlations in the decays of the fundamental particles, and also using the new simulation of tau and hadron decays the correlations in these decays which can be important in the decay of SUSY particles and the Higgs boson.

This approach was originally tested using the Randall-Sundrum model and the Minimal Supersymmetric Standard Model (MSSM). An example of the mass distribution of the quark and lepton produced in the decay $\tilde{q}_L \rightarrow q_X \tilde{\tau}_R^-$, is shown in Fig. 1. It is important that the correlations in this decay are correct as it may be possible to measure the spins of the SUSY particles using this decay mode.

An important test of our new approach for the simulation of BSM physics is the inclusion of additional models. We have therefore included the Universal Extra Dimensions model. This model is a useful “straw-man” as its particle content is similar to the MSSM but the new particles have the same spin as their Standard Model counterparts, rather than opposite spin statistics as in SUSY. The mass distribution of the quark and lepton pair in the equivalent decay chain to the SUSY chain shown in Fig. 1 is shown in Fig. 2 and compared with the analytic results for these distributions.

4 Conclusions

The Herwig++ simulation includes a number of improvements over the previous FORTRAN version for both the simulation of perturbative QCD radiation and simulation of new physics. In addition to the improvements in the simulation of the perturbative stages of the event generation process presented here improvements have been made to the simulation of QED radiation (an example of the radiation in Z decays is shown in Fig. 3) and hadron and tau decays.

The most recent version of Herwig++ is now ready for the simulation of hadron collisions and further improvements will be available in the near future.
Figure 3: The total energy of the photons radiated in Z boson decays to leptons: (a) shows the spectrum when only soft radiation is included and (b) shows the effect of including the collinear approximation for the hardest emission.

Acknowledgements

We would like to thank our collaborators on the Herwig++ project upon whose work the material presented here is based. This work was supported in part by the Science and Technology Facilities Council and the European Union Marie Curie Research Training Network MCnet under contract MRTN-CT-2006-035606.

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LATEST JET RESULTS FROM THE TEVATRON

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Recent QCD jet production measurements in \( p\bar{p} \) collisions at \( \sqrt{s} = 1.96 \text{ TeV} \) at the Tevatron Collider at Fermilab are presented. Preliminary: inclusive jet, dijet, isolated photon + jet and \( Z + \text{jets} \) measurements are compared to available perturbative QCD models.

The production of particle jets with high transverse momenta in hadronic collisions is described in perturbative Quantum Chromodynamics (pQCD) as resulting from the hard scattering of strongly interacting constituents of the colliding hadrons.

Inclusive jet rates observed in hadronic collisions at high values of transverse momenta provide a basic test of pQCD. The DØ and CDF collaborations\(^1\,^2\) have measured the inclusive jet production cross section using midpoint cone and \( k_T \) algorithms\(^3\) using data corresponding to the integrated luminosities of about 1 fb\(^{-1} \). The DØ result\(^4\) is shown in Figure 1 for two regions of rapidity\(^6\) (closed and open circles). The error bars correspond to the total measurement uncertainty. The data are corrected for the jet energy scale (JES) determined from isolated photon plus jets events, selection efficiencies and migrations due to \( p_T \) resolution (an ansatz function convoluted with the jet \( p_T \) resolution measured directly in data). The JES is the dominant source of systematic uncertainty. The integrated luminosity is known with accuracy of 6%.

The data are compared to the next-to-leading order (NLO) pQCD predictions computed using NLOJET++\(^5\) with parton density functions (PDFs) from CTEQ6.1M\(^6\) after applying threshold corrections at 2-loop (next-to-next-leading-logarithm) accuracy.\(^7\) The same jet algorithm was used in the calculations and the pQCD predictions are also corrected for hadronization effects using PYTHIA\(^8\). The theory describes the data well over the whole measured \( p_T \) range in all rapidity regions. The experimental uncertainties are competitive with those from the proton.

\(^{*}\)The rapidity \( y \) is defined as \( y = -\frac{1}{2} \ln \frac{E - p_z}{E + p_z} \), where \( E \) and \( p_z \) denote the energy and the momentum component along the proton beam direction, respectively.
PDFs and the data therefore further constrain the gluon density functions at high-\(x\). Inclusive jet spectra measured by CDF are also in good agreement with the NLO pQCD predictions.\(^9\)

The rate of dijet event produced in hadronic collisions not only provides a test of pQCD but also is sensitive to new physics such as compositeness and massive particles decays. The ratio of the dijet cross section measured by CDF to theory is shown as a function of dijet invariant mass \(M_{jj}\) in Figure 2. The measurement corresponds to 1.13 \(\text{fb}^{-1}\) and centrally produced jets. The jets were selected using midpoint cone algorithm.\(^3\) The error bars and shaded bands represent the statistical and systematic uncertainties respectively. Theoretical predictions were calculated using NLOJET++\(^5\) with PDFs from CTEQ6.1M\(^6\) and corrected to the hadron level. The systematic errors are comparable to the PDF uncertainties and NLO pQCD predictions are consistent with the data over the whole measured \(M_{jj}\) range.

The production rates of \(b\bar{b}\) jet pairs have also been studied by CDF using a data sample of 260 \(\text{pb}^{-1}\). Such a measurement provides insight into \(b\)-quark direct production, flavour excitation and gluon splitting mechanisms and also allows a test of radiative gluon corrections. The selected events were first required to contain two jets with transverse energy above 20 GeV associated to two displaced vertex tracks at the trigger level. A Run I cone algorithm\(^3\) was used to identify the jets. Jets having a positively displaced secondary with respect to the jet axis were tagged as "SVT \(b\)-jets". Two positively tagged jets with central pseudorapidities\(^8\) and transverse energies of a leading and a second jet above 35 and 32 GeV respectively were required. The invariant mass of the tracks associated to the secondary vertex was fitted using signal and background Monte Carlo templates to determine the \(b\bar{b}\) purity of the final event sample. The resulting purity was about 80%. The differential cross section measured as a function of the azimuthal angle between two jets (\(\Delta \phi_{jj}\)), unfolded to the hadron level, is shown in Figure 3 (full squares). The error bars and shaded bands correspond to the measured statistical and systematic uncertainties respectively. The data are compared to three theoretical models: two predictions at leading order (LO) from PYTHIA\(^8\) (Tune A) and HERWIG\(^10\) with PDFs from CTEQ5L\(^11\) and a NLO prediction from MC@NLO\(^12\) using CTEQ6.1M\(^6\) PDFs and with multiple parton interactions simulated by JIMMY.\(^13\) The \(\Delta \phi_{jj}\) spectrum is sensitive to contributions arising from soft gluon radiation and therefore only the NLO pQCD model describes data reasonably well even at large departures from a back-to-back jet topology.

Photons produced directly in parton-parton QCD interactions arrive unaltered at the elec-

\(^{8}\)Pseudorapidity \(\eta\) is defined as \(\eta = -\ln \tan \frac{\theta}{2}\), where \(\theta\) is the polar angle w.r.t. the proton beam direction.
tromagnetic calorimeter and carry clean information of the dynamics of the hard scatter. At the Tevatron prompt photon production is dominated by the Compton scattering subprocess $qg \rightarrow \gamma q$ for photon transverse momenta $p_T^q \lesssim 150$ GeV/c. The differential cross section $d^3\sigma/(dp_T^q d\eta^q d\eta^{\text{jet}})$ for the production of a photon and a jet measured by D0 using a 1.1 fb$^{-1}$ data sample is shown in Figure 4. The jets were reconstructed using a midpoint cone algorithm and were required to have: transverse momenta $p_T^{\text{jet}} > 15$ GeV/c and pseudorapidities either in the central calorimeter ("CC", $|\eta^{\text{jet}}| < 0.8$) or in the end cap (forward) calorimeter region ("EC", $1.5 < |\eta^{\text{jet}}| < 2.5$). Photons were selected with transverse momenta $30 < p_T^\gamma < 300$ GeV/c and central pseudorapidities $|\eta^\gamma| < 1$. Strong isolation criteria were imposed on photon candidates in order to filter background events with neutral hadrons decaying to photons in the final state. The purity of the resulting sample was estimated with the help of an artificial neural network trained to distinguish between direct photons and background. The measured cross section is corrected for the finite resolution of the calorimeter. Events with a leading jet and a photon contained in the same hemisphere in terms of their pseudorapidities are denoted as "SS" (same sign) while the remaining ones as "OS" (opposite sign). The four curves overlaid on the data represent the NLO pQCD predictions from JETPHOX with the choice of CTEQ6.1M PDFs and fragmentation functions. The theory qualitatively reproduces the data in some kinematic regions.

Jets accompanied by $W$ or $Z$ vector bosons in $p\bar{p}$ collisions constitute an important background for top quark production, Higgs and SUSY searches. Their production rates are also sensitive to physics beyond the Standard Model (compositeness and decays of heavy objects). In addition, $Z$+jets events are suitable for testing phenomenological models of the underlying event in $p\bar{p}$ collisions by studying integrated and differential jet shapes or energy flow with respect to the momentum of a $Z$ boson. The CDF collaboration studied production of $Z$+jets events with $Z$ bosons decaying into $e^+e^-$ pair using 1.1 fb$^{-1}$ of data. Such a channel provides much cleaner experimental signature than $W$+jets one, albeit has 10 times smaller cross section. The analysis covered the following kinematic region: jets reconstructed using midpoint cone algorithm having $p_T^{\text{jet}} > 30$ GeV/c and $|\eta^{\text{jet}}| < 2.1$, electrons with $E_T^{e1,2} > 25$ GeV, $|\eta^{e1}| < 1$, $|\eta^{e2}| < 2.8$ and isolated from jet cones. The acceptance window for the invariant mass of an $e^+e^-$ pair was taken to be 66 to 116 GeV/c$^2$ to suppress background. The cross section as a function of the transverse momentum of a leading jet is shown in Figure 5 (closed circles) and is compared to the NLO prediction using MCFM with CTEQ6.1M PDFs after corrections to the hadron level (open circles). The data/theory ratio is consistent with unity, although statistical errors dominate at $p_T^{\text{jet}} > 100$ GeV/c region.

Present experimental data on QCD jets from the Tevatron Collider are reasonably well
described by existing next-to-leading calculations after applying parton-to-hadron level corrections. The CDF and DØ collaborations have now collected about 2.2 fb⁻¹ of data on tape and anticipate up to 8 fb⁻¹ by end of 2009. This should make possible even higher precision tests of pQCD theory over extended kinematic regions.

Acknowledgments

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INCLUSIVE JET CROSS-SECTION MEASUREMENTS AT CDF

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Results on inclusive jet production in proton-antiproton collisions at $\sqrt{s}=1.96$ TeV based on $1fb^{-1}$ of CDF Run II data are presented. Measurements are performed using the $k_T$ algorithm in a wide range of jet transverse momentum and jet rapidity. The measured cross sections are compared to next-to-leading order perturbative QCD calculations.

The measurement of the inclusive jet cross section at the Tevatron is an important test of perturbative QCD (pQCD) predictions over more than 8 orders of magnitude, probing distances down to $10^{-19}m$. The increased center-of-mass energy in Run II (from 1.8 TeV to 1.96 TeV), the highly upgraded CDF detector, and the amount of data collected allow jet measurements in an extended region of jet transverse momentum, $p_T^{jet}$, and jet rapidity, $y^{jet}$. Jet measurements at large rapidities are important because they are sensitive to the gluon density in the proton in a kinematic region in $p_T^{jet}$ where no effect from new physics is expected.

This contribution presents results on inclusive jet production in five jet rapidity regions up to $|y^{jet}| = 2.1$, based on $1fb^{-1}$ of CDF Run II data. CDF used the longitudinally-invariant $k_T$ algorithm to search for jets:

$$k_t = p_T^{\phi,t}; \quad k_{ij} = \min(p_T^{\phi,t}, p_T^{\phi,j}) \cdot \frac{(y_i - y_j)^2 + (\phi_i - \phi_j)^2}{D^2},$$

where particles are clustered according to their relative transverse momentum. The algorithm includes a D parameter that approximately controls the size of the jet in the $\phi - y$ space. These algorithm is infrared/collinear safe to all orders in pQCD and it does not need to solve situations with overlapping jets, making possible a better comparison between data and theory. A previous
measurement using the $k_T$ algorithm at the Tevatron during Run I$^3$ observed a marginal agreement with NLO pQCD at low $p_T^{jet}$, thus suggesting the $k_T$ algorithm was particularly challenging in hadron collisions. However, these CDF results$^4$ show that this discrepancy is removed after non-perturbative corrections are included.

Figure 1 shows the measured inclusive jet cross sections using the $k_T$ algorithm with $D=0.7$, for jets with $p_T^{jet} > 54$ GeV/c in five jet rapidity regions up to $|y^{jet}| = 2.1$. For presentation, the different cross sections are scaled by a given factor. The measured cross sections have been corrected for detector effects back to the hadron level using PYTHIA-Tune A Monte Carlo$^5$, that provides an accurate description of the underlying event$^6$ and jet shapes$^7$ in Run II. The cross sections decrease over more than seven orders of magnitude as $p_T^{jet}$ increases. The systematic uncertainties on the data, mainly dominated by a 2% to 3% uncertainty in the jet energy scale, vary from 10% at low $p_T^{jet}$ to about 50% at high $p_T^{jet}$. The measurements are compared to NLO pQCD predictions as determined using JETRAD$^8$ with CTEQ6.1M PDFs$^9$ and renormalization and factorization scales set to $p_T^{max}$/2, where $p_T^{max}$ is the $p_T$ of the leading jet. The theoretical calculations include correction factors, $C_{HAD}$, to take into account non-perturbative effects related to the underlying event and fragmentation processes. The factors, presented in figure 2, have been evaluated with PYTHIA-Tune A as the ratios between the nominal cross sections at the hadron level and the ones obtained after turning off multiple parton interactions between remnants and fragmentation into hadrons. The difference obtained when HERWIG$^10$ is used instead of PYTHIA has been taken as the uncertainty on these factors. Figure 3 shows the ratios between the measurements and the theory are presented. A good agreement is observed over all $p_T^{jet}$ ranges in all rapidity regions. The uncertainty in the theoretical prediction is dominated by the uncertainty on the gluon PDF at high $x$ which, at high $p_T^{jet}$, goes from $+30\%$ to $+40\%$ for central and forward jets, respectively. The uncertainties in the data compared to that in the NLO pQCD calculations show that the measurements will contribute to a better knowledge of the parton distributions inside the proton.

![Figure 1: Inclusive jet cross sections measured using the $k_T$ algorithm with $D=0.7$ for jets with $p_T^{jet} \geq 54$ GeV/c in five rapidity regions up to $|y^{jet}| = 2.1$. The black squares represent the measured cross sections and the shaded bands indicate the total systematic uncertainty on the data. The measurements are compared to NLO pQCD calculations. The dashed lines represent the PDF uncertainties on the theoretical predictions.](image-url)
Figure 2: Parton to hadron level corrections applied to the NLO calculations to correct for underlying event and hadronization contribution in the different rapidity regions. The shaded bands represents the associated uncertainty coming from the Monte Carlo modeling.

Figure 3: Comparison between the measurements and the pQCD calculations. The dots are the ratios Data/Theory, the shaded bands indicate the total systematic uncertainty on the data and the dashed lines represent the PDF's uncertainties on the theoretical predictions.
For central jets, $0.1 < |y| < 0.7$, the measurements are repeated using a D parameter equal to 0.5 and 1.0. As D increases, the average size of the jet in $\phi - y$ space increases, and the measurement becomes more sensitive to underlying event contributions. Figure 4 shows the measurements. The good agreement still observed between the measured cross sections and the NLO pQCD predictions indicates that the soft contributions are well under control.

![Figure 4](image_url)

Figure 4: Inclusive jet cross sections measured using the $k_T$ algorithm with $D=0.5$ (left) and $D=1.0$ (right) for jets with $p_T^{jet} > 54$ GeV/c and $0.1 < |y| < 0.7$. The black squares represent the measured cross sections and the shaded bands indicate the total systematic uncertainty on the data. The measurements are compared to NLO pQCD calculations. The dashed lines represent the PDF uncertainties on the theoretical predictions. The bottom plots show the parton to hadron level corrections applied to the NLO calculations to correct for underlying event and hadronization effects, where the shaded bands represent the associated uncertainty coming from the Monte Carlo modeling.

In summary, this contribution reports results on inclusive jet production in proton-antiproton collisions at $\sqrt{s} = 1.96$ TeV, based on $1fb^{-1}$ of CDF Run II data, using the $k_T$ algorithm. CDF also performed the measurement using the Midpoint cone-based algorithm. The measurements are in good agreement with NLO pQCD calculations. In particular, for central jets and at high $p_T^{jet}$ no deviation with respect to the theory is found. In the most forward region, the total systematic uncertainty on the data is smaller than that on the theoretical calculations. Therefore, these new results will contribute to a better understanding of the gluon PDF in the proton at high $x$.

References

Exclusive $e^+e^-$, Di-photon and Di-jet Production at the Tevatron

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Results from studies on exclusive production of electron-positron pair, di-photon, and di-jet production at CDF in proton-antiproton collisions at the Fermilab Tevatron are presented. The first observation and cross section measurements of exclusive $e^+e^-$ and di-jet production in hadron-hadron collisions are emphasized.

1 Introduction

Exclusive production in $\bar{p}p$ collisions is a process in which the proton and antiproton remain intact and an exclusive state $X_{excl}$ of particle(s) is centrally produced: $\bar{p} + p \rightarrow \gamma \gamma + X_{excl} + p'$. Exclusive di-lepton production may occur through a two photon exchange process; $\bar{p} + p \rightarrow \gamma \gamma \rightarrow e^+e^-$, where the lepton is an electron. Since this process, shown in Fig. 1 (left), depends essentially only on Quantum Electro-Dynamics, the cross section is known with an accuracy better than 1 %. Exclusive di-jet (or di-photon) production is a special case of di-jet (di-photon) production in double Pomeron exchange (DPE), a diffractive process in which the $p$ and $\bar{p}$ remain intact with a small momentum loss, and a system $X_{incl}$ containing the jets (photons) is produced: $\bar{p} + p \rightarrow [p' + IP_p] + [p' + IP_p] \rightarrow \gamma \gamma + X_{incl} + p'$. Here, $IP$ represents a Pomeron, defined as an exchange consisting of a colorless combination of gluons and/or quarks carrying the quantum numbers of the vacuum. In a particle-like Pomeron picture (e.g. see 1), the system $X_{incl}$ produced by two Pomeron collisions, $IP_p + IP_p \rightarrow X_{incl} \Rightarrow (jet_1 + jet_2) + Y$, generally contains Pomeron remnants $Y$ in addition to jets (in di-jet case). Di-jet or di-photon production by DPE may occur as an exclusive process with no Pomeron remnants ($Y = 0$). Exclusive di-jet or di-photon production may also occur through a $t$-channel color-singlet two gluon exchange at leading order (LO) perturbative Quantum Chromo-Dynamics (QCD), shown in Fig. 1 (center). A similar diagram (right) is used to calculate exclusive Higgs Boson production 2.

In this paper, we report an observation of exclusive $e^+e^-$ production by CDF in hadron-hadron collisions 3. Also, reported is a first observation of exclusive di-jet production in $\bar{p}p$
collisions, and preliminary results on exclusive di-photon search at CDF.

2 CDF II Detector

Among relevant CDF II detectors, crucial for detecting the rapidity gap (defined as a region of pseudorapidity devoid of particles) in the exclusive events are the forward detectors which consist of the MiniPlug calorimeters (MPCAL)\(^4\), the Beam Shower Counters (BSC), a Roman Pot Spectrometer (RPS), and Cerenkov Luminosity Counters (CLC)\(^5\). The MPCAL, designed to measure particle energies in the region \(3.6 < |\eta| < 5.2\), consist of alternating lead plates and liquid scintillator layers perpendicular to the beam, which are read out by wave-length shifting fibers. Combining with the central (CCAL) and plug (PCAL) calorimeters, the MPCAL extends the calorimeter coverage up to \(|\eta| = 5.2\). The BSC are scintillation counters surrounding the beam pipe at three (four) different locations on the outgoing \(p\) (\(\bar{p}\)) side of the CDF II detector. Covering the range \(5.4 < |\eta| < 5.9\) is the BSC1 system, which is used for triggering on events with forward rapidity gaps. The RPS, located at \(\sim 57\) m downstream in the \(p\) beam direction, consists of three Roman pot stations, each containing a scintillation counter used for triggering on the \(\bar{p}\), and a scintillation fiber tracking detector for measuring the position and angle of the detected \(\bar{p}\). The CLC covering the range \(3.7 < |\eta| < 4.7\) are used to measure the luminosity and, in some cases, to sharpen the rapidity gap in the MPCAL.

3 Data Samples

Main data samples used in the studies presented in this paper are collected with two dedicated diffractive triggers. Search for exclusive \(e^+e^-\) and di-photon production uses the data triggered on two clusters of energy in the CCAL or PCAL calorimeters, and no activity in the BSC1 counters on both sides. Off-line, events containing electron or photon candidates with \(E_T > 5\) GeV are selected. The total integrated luminosity of this sample is \(532 \pm 32\) pb\(^{-1}\).

Exclusive di-jet search is performed on a data sample triggered when the event contains 1) a 3-fold coincidence of the RPS trigger counters in time with a \(\bar{p}\) gate, b) at least one trigger tower with \(E_T > 5\) GeV, and c) no BSC1 activity on the outgoing \(p\) side (referred to as RP+GAP+Jet5 trigger). The events in the sample with an integrated luminosity of \(313 \pm 19\) pb\(^{-1}\) are required to pass the following selection cuts; 1) no more than one reconstructed primary vertex within \(|z| < 60\) cm, 2) RPS trigger counter pulse height cut, 3) at least two jets (corrected to hadron level) with \(E_T > 10\) GeV within \(|\eta| < 2.5\), 4) \(0.01 < \xi^X_\rho < 0.12\), and 5) zero MPCAL and zero CLC hit multiplicity on the \(p\) side. Cut 1) is used to reduce overlap events with multiple \(pp\) collisions occurring during the same bunch crossing. Cut 2) is imposed to reject triggers caused by particles hitting the beam pipe in the vicinity of the RPS. The \(\xi^X_\rho\), defined as \(\xi^X_\rho = \sum_{i=1}^{N_{\text{jet}}}(E_T^i e^{-\eta \rho})/\sqrt{E_T^i}\) where the sum is carried over all calorimeter towers with \(E_T > 100\) MeV for CCAL and PCAL, and \(E_T > 20\) MeV for MPCAL, representing the fractional momentum loss of the \(\rho\), is required to be in the range of cut 4) to reduce remaining
overlap events. Finally, the DPE contribution in the sample is enhanced by the cut 5) which requires a wide p-side rapidity gap in the region of $3.6 < \eta < 5.9$.

4 Exclusive Electron-Positron Production

Search for exclusive $e^+e^-$ signal is based on selecting events with electron and position candidates and refining the selected sample by requiring no additional particle signatures in the calorimeters or BSC. The electron (or position) candidate is defined as a cluster in the EM calorimeter with $E_T > 5$ GeV and $|\eta| < 2$, being consistent with an $e^-$ ($e^+$), and a track with $p_T > 1$ GeV/c pointing to the calorimeter cluster. A particle signature in the calorimeter (BSC) is defined as a cluster of adjacent towers in the MPCAL or a single tower in the CCAL or PCAL (any BSC hit) above the noise threshold. From the sample of triggered events, a total of 16 candidate events pass electron identification, cosmic ray rejection, and the veto on additional particle signatures.

The total background, which consists of a) jets faking electrons, b) cosmic rays interacting in the detector, c) non-exclusive events, and b) $\gamma\gamma \rightarrow e^+e^-$ events with proton dissociation, is estimated to be $1.9 \pm 0.3$ events in the 16 candidate events. The dominant contribution is proton dissociation events ($1.6 \pm 0.3$ events). The cross section for exclusive $\bar{p}p \rightarrow \bar{p} + e^+e^- + p$ is measured to be $1.6^{+0.8}_{-0.3} (\text{stat}) \pm 0.3 (\text{syst})$ pb, which is in good agreement with the theoretical cross section of $1.71 \pm 0.01$ pb obtained from LPAIR Monte Carlo (MC) simulation. The probability of observing 16 events or more when $1.9 \pm 0.3$ events are expected is $1.3 \times 10^{-9}$, equivalent to a $5.5\sigma$ effect. The good agreement between the measured and expected cross sections proves that rapidity gap signature in the detectors is well understood.

5 Exclusive Di-photon Production

Search for exclusive di-photon event, $\bar{p}p \rightarrow \bar{p} + \gamma\gamma + p$, follows the same event selections as the exclusive $e^+e^-$ search, except that the photon candidates (defined as EM calorimeter clusters with $E_T > 5$ GeV and $|\eta| < 1$) have no tracks pointing to the clusters. The total event selection efficiency for exclusive $\gamma\gamma$ events with photons of $E_T > 5$ GeV and $|\eta| < 1$ is $4.0 \pm 0.7$ %.

In the same sample as used in exclusive $e^+e^-$ analysis, three events passed the selection criteria. The ExHuME MC generator for exclusive di-photon production via $gg \rightarrow \gamma\gamma$ processes predicts $1^{+3}_{-1}$ events, which is consistent with the number of observed candidate events. Background estimation is currently under way.

6 Exclusive Di-jet Production

A special analysis strategy is developed to search for the exclusive di-jet production as it is quite difficult to identify particles associated with the jets. The analysis employs distributions of "di-jet mass fraction", $R_{jj}$, defined as the di-jet invariant mass $M_{jj}$ divided by the whole system mass $M_X$ (except leading nucleons); $R_{jj} = M_{jj}/M_X$. The $M_X$ is obtained from all calorimeter towers above the thresholds used to calculate the $e^\pm_X$, and the $M_{jj}$ is calculated from calorimeter tower energies inside the $R=0.7$ cones of jets. The exclusive signal is extracted by comparing the DPE di-jet data selected in Sec. 3 with inclusive DPE di-jet events (which do not contain exclusive di-jet production) in the $R_{jj}$ distribution shape. We use POMWIG MC generator with detector simulation to simulate the DPE di-jets.

The $R_{jj}$ distribution comparisons showed a clear excess of data at high $R_{jj}$, as expected for exclusive di-jet signal. It turns out that the observed excess is robust against the choice of Pomeron parton distribution functions and Pomeron remnant effects on underlying events. In order to extract the signal from the data, two exclusive di-jet MC simulations, ExHuME and exclusive DPE model (EXCLDPE) in DPEMC have been studied. Fig. 2 (left) shows the fit
Figure 2: Left: Dijet mass fraction in DPE data (points) and best fit (solid histogram) to the data obtained from POMWIG MC events (dashed histogram) and ExHuME di-jet MC events (shaded histogram). Center: Measured exclusive di-jet cross section for $R_{jj} > 0.8$ (points) as a function of minimum second jet $E_T$, compared with predictions from ExclDPE (dashed curve) and ExHuME (dotted curve) MC simulations, and the KMR calculation at LO parton level, scaled down by a factor 3 (shaded band). Right: Values of $F_1$ (filled points) and $F_2$ (open points) as a function of $R_{jj}$, where $F_1$ is the ratio of heavy flavor jets to all inclusive jets, normalized to the weighted average value in the region of $R_{jj} < 0.4$ (systematic uncertainties are shown by the shaded band), and $F_2$ is the ratio of POMWIG MC to inclusive DPE di-jet data.

to the data $R_{jj}$ shape using a combination of POMWIG inclusive DPE di-jets and ExHuME exclusive di-jets for events with di-jets of $E_T > 10$ GeV and the third jet veto of $E_T^{jet} < 5$ GeV. The third jet veto is introduced because the exclusive MC generates only LO $gg \rightarrow gg$ process. The fit shows the data excess is well described by the presence of exclusive di-jets. The EXCLDPE provides similar results in the shape fit. From the MC fits to the data, we measure the cross section of exclusive di-jet production as a function of minimum second jet $E_T$, as shown in Fig. 2 (center). The data prefer ExHuME and perturbative QCD calculations at LO parton level (KMR) in Ref. 2.

Exclusive $gg \rightarrow q\bar{q}$ production is expected to be suppressed at LO by the $J_z = 0$ total angular momentum selection rule. Confirming the suppression will therefore support the results obtained from the MC based method described above. Since the suppression holds for any quark flavor, we measure the ratio $F_1$ of heavy flavor ($b$ or $c$) quark jets to all jets as a function of $R_{jj}$ using a $200$ pb$^{-1}$ sample of data triggered on RP+GAP+Jet5 and a transversely displaced track of $p_T > 2$ GeV. The result, presented in Fig. 2 (right), shows $F_1$ versus $R_{jj}$, where the $F_1$ is normalized by the weighted average value in the range of $R_{jj} < 0.4$. The decreasing trend of $F_1$ with $R_{jj}$, which indicates a manifestation of the $J_z = 0$ selection rule, is compared with the MC based result given as $F_2$, where $F_2$ is the ratio of the inclusive MC predicted events (normalized to the data at $R_{jj} < 0.4$) to the data. The two results are consistent with each other in both magnitude and $R_{jj}$ dependence.

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A Practical Seedless Infrared Safe Cone Algorithm

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This writeup highlights the infrared unsafety of the 'midpoint' cone jet-algorithm and provides a brief overview of why this is a serious issue. It then shows how one can build a safe (seedless) cone algorithm and discusses the potential impact on measurements.

Two broad classes of jet algorithm are in widespread use at modern colliders: sequential recombination type algorithms, such as $k_t^4$ and Cambridge/Aachen$^2$, and cone-type ones$^3$. The former take a bottom-up approach to the problem of defining jets, repeatedly combining particles that are closest in some distance measure. They work because the proximity measures used are closely related with QCD divergences for particle production, and they are much appreciated in the $e^+e^-$ and $ep$ communities, both for their simplicity and their modest hadronisation corrections. Cone type algorithms take a top-down approach, finding coarse regions of energy flow (cones) and identifying them as jets. They work because QCD only modifies the energy flow on small scales, and so far they have been the preferred type of algorithm in the $pp$ community, because of the greater geometrical regularity of the resulting jets and their sometimes lower sensitivity to certain components of the non-perturbative underlying event and pileup.

Cone algorithms have been in use since the early 1980's$^4$, and in the early 1990's awareness developed$^5$ of the importance for cone-algorithm formulations to satisfy a certain basic set of requirements: they must be fully defined, practical in both experimental and theoretical contexts, and cross sections must be finite at any order of perturbation theory, i.e. the algorithm must be infrared and collinear (IRC) safe.

Modern cone algorithms involve two main steps: a procedure to find ‘stable cones’ (a cone pointing in the same direction as the momentum of its contents) and a ‘split–merge’ procedure to convert those cones into jets, resolving the problem of cones that have particles in common (i.e. that ‘overlap’).

$^a$Talk presented at the XLII Rencontres de Moriond, QCD and Hadronic Interactions, La Thuile, March 2007.
The most delicate issue with cone algorithms has been that of finding the stable cones. A standard procedure had long been to use all particles (possibly above a seed threshold) as directions of trial cones, then for each trial cone to use the momentum of its contents as a new trial direction, iterating until stable directions are obtained. The drawback of iterative procedures is that new stable cones (and jets) may be found if an extra starting point is added. This was known to happen with the addition of soft particles in straightforward iterative stable-cone searches, but it had been thought that a trick of adding extra starting points, at the midpoints between the cones already found, would lead to a final set of stable cones that was insensitive to the addition of extra seeds. Accordingly a recommendation was made for the Tevatron experiments to use such a ‘midpoint’ iterative cone algorithm.

It turns out that while the midpoint fix resolves the infrared problems for events with two neighbouring hard particles, those problems reappear for three or more neighbouring hard particles. This is illustrated in fig. 1 where in the left-hand particle configuration two stable cones (and so two jets) are found with the midpoint cone algorithm. If a soft, \( \sim 1 \text{ GeV} \), particle is now added (right), it provides an extra seed leading to a third (overlapping) stable cone being found, and all the cones are then merged into a single jet (for \( f = 0.5 \)).

The sensitivity to the set of seeds means that the midpoint cone algorithm is either infrared unsafe (without a seed threshold) or collinear unsafe (with a seed threshold). This is a serious issue, for many reasons: 1) it defeats the purpose of using a jet algorithm in the first place: a jet algorithm is supposed to provide a correspondence between the complex hadron level and a simple few-parton picture of an event — this correspondence is meaningless if a random 1 GeV non-perturbative particle changes the multi-hundred GeV jets. 2) IRC unsafety invalidates the theorems that ensure the finiteness of perturbative QCD calculations, because the jets found in (divergent, supposedly cancelling) real and virtual diagrams differ. 3) Pragmatically it limits the accuracy with which one can meaningfully predict many observables, as summarised in table 1, and already programs such as NLOJET or MCFM allow one to go beyond this order when using a safe jet algorithm. Therefore the use of a midpoint algorithm squanders the potential for accurate predictions that stems from many years of hard theoretical calculations, and forever limits the usefulness of data measured with it.

A solution to the cone algorithm’s problems is to carry out an exhaustive (‘seedless’) search for all stable cones. Since additional soft particles do not change the stability of cones, if one

<table>
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<tr>
<th>Observable</th>
<th>1st miss cones at</th>
<th>Last meaningful order</th>
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<td>NNLO</td>
<td>NLO</td>
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<td>( W/Z/H + 1 ) jet cross section</td>
<td>NNLO</td>
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<td>3 jet cross section</td>
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<tr>
<td>( W/Z/H + 2 ) jet cross section</td>
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<td>jet masses in 3 jets, ( W/Z/H + 2 ) jets</td>
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Table 1: Summary of the order \( (\alpha_s \text{ or } \alpha_s^2 \log W) \) at which stable cones are missed for various observables with a midpoint algorithm, and the corresponding last order that can be meaningfully calculated. (Legacy iterative cone algorithms, without midpoint seeds, such as JetClu, fail one order earlier).
has already found all stable cones adding a soft particle cannot lead to extra stable cones being found, and so the IRC safety problem is eliminated. A seedless algorithm had been proposed\textsuperscript{11,13} for perturbative calculations, but since it took time $\sim N2^N$ to find jets among $N$ particles ($10^{17}$ years for $N = 100$), it was unthinkable to use it at hadron or detector level.

Recently it was observed\textsuperscript{12} that it can be advantageous to relate sequential-recombination jet algorithms to problems in computational geometry. It turns out that this is true also of cone algorithms, for which the exhaustive stable cone search reduces to a 2D ‘all distinct circular enclosures’ problem. While apparently not having been considered by the computational geometry community, this problem is easily solved, essentially by considering all circles having a pair of particles on their circumference, cf. fig. 2. With the aid of further standard computational techniques one obtains\textsuperscript{8} a seedless algorithm that takes $\mathcal{O}(N^2 \ln N)$ time. Not only does this provide a practically usable IR safe cone algorithm, but it even scales better at large $N$ than midpoint algorithms ($N^3$) and is of similar speed to them for the values of $N \sim 500 - 1000$ that will be found at low-luminosity LHC.

Given the cone algorithm’s chequered history with IRC safety, it is important to establish, as far as possible, that there are no further unpleasant surprises waiting to be discovered in a few years’ time. This has been done in two ways: with a detailed analytical proof, and via Monte Carlo tests in which one finds jets in a ‘hard event’ (with up to 10 hard particles), repeatedly adds infinitely soft particles and verifies that the jets found are the same. If they are not, then the algorithm is IR unsafe. The failure rates on this test are shown for a variety of cone algorithms in fig. 3. Among the discoveries made in these tests, was that the split-merge procedure also had the potential to create IR safety problems. The final version of the seedless algorithm, named SISCon, has passed several billion hard event tests without failure. The code for the algorithm is available publicly\textsuperscript{13} both in standalone form and as a FastJet\textsuperscript{12} plugin.

The physics impact of switching from the midpoint to SISCon depends on the observable and is illustrated in fig. 4 for two cases. For inclusive quantities, like the inclusive jet spectrum (upper panel), one sees effects of the order of a couple of percent, as is to be expected since...
stable cones are only missed at NNLO onwards. One notes nevertheless that differences of up to 5% arise when including underlying event effects, and this is related to SISCone’s substantially lower sensitivity to diffuse ‘noise’ in an event.

For more exclusive quantities the differences between midpoint and SISCone are more significant. For jet-mass spectra in three-jet events (lower panel of fig. 4), the difference starts are LO, and this can translate to 40% effects in partonic predictions (which essentially corresponds to an unavoidable 40% non-perturbative ambiguity for the midpoint algorithm).

To conclude, while both sequential recombination and cone-type jet algorithms have their place at hadron colliders, it is essential that they be practical and safely defined. The widespread ‘midpoint’ cone algorithm is not infrared safe, and therefore there are strong reasons for discontinuing its use in favour of a seedless cone algorithm such as SISCone, which is both infrared safe and practical at parton, hadron and detector levels.

Acknowledgements

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Jet Areas, and What They are Good For

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We introduce the concept of the area of a jet, and show how it can be used to perform the subtraction of even a large amount of diffuse noise from hard jets.

1 Introduction

Jet clustering algorithms, which map the particles observed in the final state of a high-energy collisions into a smaller number of (usually) well defined objects – the jets – are widely used in the study of the properties of strong interactions. The jets are usually meant to be good proxies of the original partons (though the detailed relation is more subtle), and by studying them one tries to probe the underlying dynamics. The reason for using the jets, rather than directly the observed hadrons, is that they can be construed as infrared-safe observables: they are therefore amenable to perturbative QCD predictions, and their sensitivity to non-perturbative phenomena (hadronisation, underlying event and pileup effects) can either be kept under control or corrected for.

In this talk we explore the issue of the susceptibility of jets to contamination from soft radiation distributed in the form of a roughly uniform and diffuse background. Physical examples are the pileup originated by multiple minimum bias collisions in high-luminosity hadron colliders like the LHC, the many particles produced in a central heavy ion collision and, to a lesser extent, the underlying event given by perturbative and non-perturbative QCD radiation whenever strongly-interacting particles are produced at high energy. We shall argue that this susceptibility can be quantitatively characterised in terms of the novel concept of area of a jet, which we shall rigorously introduce. In turn, this will suggest a procedure by means of which such contamination can be subtracted from the jet momentum, so as to recover – to a large extent – its proxy relation with the parton it originated from.

*In collaboration with Gavin Salam and Gregory Soyez.
Figure 1: a) Active area distributions for the $k_t$ algorithm. Cambridge/Aachen has a very similar behaviour.

Naively, one can think of the jet area as the surface (in the rapidity-azimuth plane) over which the particles that have been clustered into a given jet are distributed. One can also assume that the amount of diffuse background radiation clustered together with the jet will be proportional to this area. One could therefore think of determining somehow the momentum surface density of this noise, $\rho$, and successively subtract from the jet momentum a quantity given by $\rho$ times the area of the jet.

Before such a program can be implemented in practice, however, the jet area needs to be defined more rigorously, and a procedure to extract $\rho$ must be devised. This is done in and respectively, where both aspects are introduced and extensively studied.

2 Jet Area

The naive vision of the jet area as the surface covered by the particles that make up the jet quickly turns out to be fallacious: as the particles are point-like, this area is zero. Drawing some sort of boundary, like for instance a convex hull – the minimal set of particles such that all the others are contained in the polygon drawn through them – is also prone to ambiguities: different jets may overlap, and a region of space might be arbitrarily assigned to a jet irrespectively of the properties of the clustering algorithm.

To overcome these difficulties, we propose a definition of jet area which is inherently related to the clustering procedure, and which can properly account for the jet contamination due to a diffuse background. Our definition is strictly dependent on the infrared-safety property that a good jet algorithm should have: the addition of one (or many) soft particles to the event should not change the final set of hard jets. We add therefore a large number of uniformly distributed and extremely soft particles (ghosts) to the event, and cluster them together with the real particles. At the end of the clustering procedure, the number of ghosts clustered with each jet will provide a robust measure of the jet’s extension in the rapidity-azimuth plane, and define therefore its active area, $A_h$.

Fig. ??(a) shows how the values for this active area are distributed for two kinds of events: on one extreme, jets constituted of many uniformly distributed particles with similar momenta (the pure-ghost jets); on the other extreme, a jet containing a single hard particle. We can see that these two situations produce different distributions for the active areas, with different averages:

---

4 The drawback of this procedure is that a very large number of particles needs to be clustered (a few thousands ghosts are needed to achieve accuracies of the order of one per cent). This would be unfeasible – or at least extremely impractical – without the fast implementations of the $k_t$ and the Cambridge/Aachen jet algorithms provided by FastJet. This package also provides the tools to calculate the area of the jets, as well as an interface to the new infrared-safe cone algorithm SISCone.
the jets containing many similar particles have a typical area of order \( \langle A_{\text{soft}} \rangle \simeq 0.55 \pi R^2 \) (\( R \) is the typical radius parameter present in most jet algorithms), while the jets containing a single hard particle tend to be larger, their average area being \( \langle A_{\text{hard}} \rangle \simeq 0.81 \pi R^2 \).

One can take further this exploration of similarities and differences between soft (i.e. uniform) and hard jets, and explore how the transition takes place: \( \text{fig. ??(b)} \) shows the average area of the jet containing the single "hard" particle as its transverse momentum \( p_t \) changes from being negligible with respect to the soft background to being much larger. One can see that in the \( p_{t, \text{hard}} \gg (p_{t, \text{soft jet}}) \) limit the \( \simeq 0.81 \pi R^2 \) value for the average area is recovered. On the other hand, in the opposite \( p_{t, \text{hard}} \ll (p_{t, \text{soft jet}}) \) limit the "hard" jet now behaves like a soft one, the difference in average area being only of probabilistic nature related to the "measurement" of the area of the specific jet containing a given particle.

3 Noise Level

The estimation of \( \rho \), the typical level of the background radiation, could probably be performed in many ways. The method we propose here is related to the jet areas discussed above. It relies on the observation that the transverse momentum of a jet divided by its area, \( p_{t,j} / A_j \), behaves differently for the hard jets and for the background ones. Typically, the jets originating from the background radiation cluster themselves in a band, while the hard jets stick out. This is clearly shown in \( \text{fig. ??} \). This event is a simulated \( pp \) collision at the LHC at moderate luminosity: 10 additional minimum bias events are added to the main hard collision, which produces a dijet event with jets of transverse momentum of the order of 50 GeV. \( \text{Fig. ?? (left)} \) shows that the areas of the various jets can fluctuate widely. However, when the same jets are plotted in terms of \( p_{t,j} / A_j \) (right plot) one clearly see the band established by the background. Different strategies can be devised to quantitatively determine its level. One of the simplest one is to take the median of all the \( p_{t,j} / A_j \), an operation that prevents the few hard jets from biasing its value.

We define therefore:

\[
\rho = \text{median} \left[ \left\{ \frac{p_{t,j}}{A_j} \right\} \right].
\]

In the specific case of the event of \( \text{fig. ??} \), the momentum density of the background is therefore \( \rho \simeq 6 \text{ GeV per unit area} \).

4 Background Subtraction

Once the area of each jet, \( A_j \), and the noise level \( \rho \) are known, one can correct the transverse momentum via the following operation:

\[
p_{t,j}^{(\text{sub})} = p_{t,j} - \rho A_j.
\]
We show how this works in practice by considering the following toy model: we generate many events which contain a single hard particle, with a transverse momentum $p_T^{hard} = 100$ GeV, embedded in a background of 10000 soft particles, each with an average transverse momentum $\langle p_T^{soft} \rangle = 1$ GeV (with little fluctuations, 10%, around this value) and randomly uniformly distributed in rapidity and azimuth up to $y_{max} = 4$. In this particular case we can of course calculate the transverse momentum density (per unit area) of the soft particles from the input parameters, since we know how we generated them:

$$\rho = \frac{\langle dp_T^{soft} / dy d\phi \rangle}{2 \times y_{max} \times 2\pi} = \frac{10000 \times 1 \text{ GeV}}{2 \times 4 \times 2\pi} \approx 200 \text{ GeV}.$$  \hspace{1cm} (3)

This situation might look extreme, but similar values are expected in realistic cases, like a central Pb Pb collision at the LHC.

![Figure 3: Jets containing a hard particle with $p_T^{hard} = 100$ GeV clustered together with a soft background (green, “raw” histogram), and after its subtraction (blue “corrected” one). The ‘4-vector’ versions of the area and of the subtraction \cite{catani}, more appropriate for large $R$, have been used for this plot.]

We know from the previous section that an average soft jet, when clustered with the $k_T$ or the Cambridge/Aachen algorithm with $R = 1$, has an area of order $0.55\pi$. This translates in a typical transverse momentum $p_T^{soft \text{ jet}} \approx \rho(A^{soft}) \approx 350 \text{ GeV}$. Such jets would already dwarf the hard particle of 100 GeV. However, this particle will itself be embedded in a jet containing also many soft particles: this jet will therefore have a typical transverse momentum of the order of $350 + 100$ GeV, but huge fluctuations will be visible from one event to another, as the amount of background clustered with it will vary considerably.

This means that both the absolute energy scale and the energy resolution are degraded by the presence of the background, as shown in fig. 3: the transverse momentum of the hard jet is displaced, by an amount consistent with our estimate, and the resolution is hopelessly bad (green histogram, “raw”). However, once the subtraction is performed according to eq. (3) (using for each event the $\rho$ directly extracted from the clustering, as explained in Sec. 3, and not the fixed value of eq. (2), of course), the correct average transverse momentum is recovered, together with a large fraction of the resolution (blue histogram, “corrected”).

This toy model shows the feasibility and the accuracy of the determination of the noise level and of the subtraction procedure. More realistic examples, and references to experimental investigations of the problem of background subtraction, are given in \cite{cacciari2}.

Acknowledgements. This work has been performed, and is partially in progress, with Gavin Salam and Gregory Soyez, whom I thank for an extremely stimulating collaboration.

References

1. M. Cacciari, G.P. Salam and G. Soyez, LPTHE-07-02, in preparation
Accurate predictions for heavy quark jets

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Heavy-flavour jets enter many of today's collider studies, yet NLO predictions for these quantities are subject to large uncertainties, larger than the corresponding experimental errors. We propose a new, infrared safe definition of heavy-quark jets which allows one to reduce theoretical uncertainties by a factor of three.

1 Introduction

When looking at the current comparison between the inclusive $b$-jet spectra measured by CDF and the corresponding next-to-leading order (NLO) predictions, Fig. 1, one notices two striking features. Firstly, one sees a tension between data and theory: the ratio of data over NLO is around 1.2-1.5 over the whole range of accessible transverse momenta $p_T$ of the jets. Secondly, one notices that the uncertainties associated with the theoretical predictions are embarrassingly

![Figure 1: Ratio of the measured inclusive $b$-jet spectrum to the NLO prediction. The measurement is performed for jets with transverse momentum $38$ GeV < $P_T$ < $400$ GeV and rapidity $|y_{jet}| < 0.7.$]
large ($\sim 40 - 50\%$) for a NLO calculation and in particular they are larger than the corresponding experimental uncertainties. To understand why this happens it is useful to examine Fig. 2.

The top plots show that the large uncertainty is associated with very large $K$-factors. The middle plots confirm that the uncertainty is the same both with MCFM$^2$ and MCNLO$^3$. Finally, the bottom plots illustrate the origin of the poor convergence of the perturbative expansion: when breaking down the Herwig$^4$ b-jet spectrum into the hard underlying channels it turns out that two NLO channels, flavour excitation, where a $b$-quark is kicked out of the sea-quarks, and gluon splitting, where a gluon in the final state splits into a $b\bar{b}$-pair, are larger than the leading order heavy quark production mechanism, flavour creation, when two incoming light partons produce a heavy quark pair.

The reason why supposedly higher order contributions are actually larger than the leading order channel can be clarified by counting soft and collinear logarithms associated with the splitting of gluons into $b\bar{b}$-pairs. It turns out that flavour excitation contributes with $(a_s \ln p_t/m_b)^n$ and gluon splitting contributes with $(a_s \ln p_t/m_b)^{2n-1}$ relative to the leading order, $O(a_s^2)$ process. Since $m_b < p_t$ these contributions are enhanced. Moreover, the dominant contribution to the b-jet spectrum comes from jets originated from gluon splitting, which do not correspond to one’s intuitive physical idea of a b-jet, one where a hard b is produced directly in the hard scattering. In the following we suggest to adopt a different jet-clustering algorithm to reconstruct b-jets. One that by making explicit use of the flavour information eliminates all higher-order logarithmic enhancements associated to gluon splittings in the b-jet spectra. This means that, after resumming initial state collinear logarithms into b-pdfs, b-jets can be computed using massless QCD calculations$^5$ as long as one neglects power corrections $m_b^2/p_t^2$ (potentially log-enhanced).

2 The heavy-quark jet algorithm

We summarize here the inclusive heavy-flavour jet algorithm for hadron-hadron collisions$^6$. For any pair of final-state particles $i$, $j$ define a class of distances $d_{ij}^{(F,a)}$ parametrized by $0 < \alpha \leq 2$ and a jet radius $R$

$$d_{ij}^{(F,a)} = \frac{R_{ij}^2}{R^2} \times \begin{cases} \max(k_{t1}, k_{t2})^\alpha \min(k_{t1}, k_{t2})^2 - \alpha, & \text{softer of } i, j \text{ flavoured,} \\ \min(k_{t1}^2, k_{t2}^2), & \text{softer of } i, j \text{ flavourless,} \end{cases}$$

\(^aWe recall that according to the current experimental definition of a b-jets, a b-jet is any jet containing at least one b.
where $R_{ij}^2 = \Delta y_{ij}^2 + \Delta \phi_{ij}^2$, $\Delta y_{ij} = y_i - y_j$, $\Delta \phi_{ij} = \phi_i - \phi_j$ and $k_{ti}$, $y_t$ and $\phi_t$ are respectively the transverse momentum, rapidity and azimuth of particle $i$, with respect to the beam. For each particle $i$ define a distance with respect to the beam $B$ at positive rapidity,

$$d_{iB}^{F,\alpha} = \begin{cases} \max(k_{ti}, k_{B}(y_i))^{\alpha} \min(k_{ti}, k_{B}(y_i))^{2-\alpha}, & i \text{ is flavoured,} \\ \min(k_{ti}, k_{B}^{2}(y_i)), & i \text{ is flavourless,} \end{cases}$$

with

$$k_{B}(y) = \sum_i k_{ti} \left( \Theta(y_i - y) + \Theta(y - y_i) e^{y_i - y} \right).$$

Similarly define a distance to the beam $B$ at negative rapidity by replacing $k_{B}$ in eq. (2) with $k_{B}^{\bar{B}}$.

$$k_{\bar{B}}(y) = \sum_i k_{ti} \left( \Theta(y - y_i) + \Theta(y_i - y) e^{y - y_i} \right).$$

Identify the smallest of the distance measures. If it is a $d_{ij}^{F,\alpha}$, recombine $i$ and $j$ into a new particle, summing their flavours and 4-momenta; if it is a $d_{iB}^{F,\alpha}$ (or $d_{i\bar{B}}^{F,\alpha}$) declare it to be a jet and remove it from the list of particles. Repeat the procedure until no particles are left. We define the $b$-flavour or generally the heavy-flavour of a (pseudo)-particle or a jet as its net heavy flavour content, i.e. the total number of heavy quarks minus heavy anti-quarks.

The IR-safety of this algorithm was proved in $^6$. Apart from allowing one to take the limit $m_Q^2 \to 0$ for the heavy quark mass (as long as collinear singularities associated with incoming heavy quarks are factorized into a heavy quark PDF), it ensures that one obtains the same results whether one considers heavy-quark flavour at parton level, or heavy-meson flavour at hadron level, modulo corrections suppressed by powers of $\Lambda_{QCD}/p_t$.

3 Results

Our results are summarized in fig. 3 $^7$ where we show the inclusive $b$-jet $p_t$-spectrum as obtained with the flavour algorithm specified above with $\alpha = 1$, and $R = 0.7$, the latter having been shown to limit corrections associated with the non-perturbative underlying event $^8$. The left (right) column of the figure shows results for the Tevatron Run II (LHC). We have selected only those jets with rapidity $|y| < 0.7$. We also show the full inclusive jet spectrum (all jets) as obtained with a standard inclusive $k_t$-algorithm $^9$ with $R = 0.7$.

We notice the considerable reduction of $K$-factors, which are around 1.3 and the moderate uncertainties associated with scale variation, signaling that the perturbative expansion is now well under control. Our predictions constitute therefore the first accurate predictions for inclusive heavy quark jets.

We remark that very similar results are obtained when considering charmed jet spectra. An interesting issue there is that predictions are very sensitive to possible intrinsic charm components of the proton $^{10}$. This means that this type of observable has a potential to set constraints on such intrinsic components.

A last remark concerns the feasibility of the experimental measurement of heavy flavour jets defined with our flavour algorithm. Our jet-clustering algorithm requires that one identify heavy-flavoured particles and that one uses a different distance measure when clustering heavy or light objects according to eq. (1). It is particularly important to identify cases when both heavy flavoured particles are in the same jet, so as to label this jet a gluon jet and eliminate it from the $b$-jet spectrum. Experimentally techniques for double $b$-tagging in the same jet already exist $^{11}$ and steady progress is to be expected in the near future $^{12}$. However one has always a limited efficiency for single $b$ tagging, and even more for double $b$-tagging in the same jet. On the other hand preliminary studies indicate that one does not necessarily need high efficiencies,
Figure 3: Inclusive jet spectrum at the Tevatron (right) and at the LHC (left). The top two panels show results for both $b$-jets and all-jets, while the lower three panels apply only to $b$-jets. See text for further details.

but what is more crucial is that one understand these efficiencies well. We look forward to further investigation in this direction.

Acknowledgments

This work is done in collaboration with Andrea Banfi and Gavin Salam.

1. CDF Collaboration, Note 8418.
Recent results of $J/\psi$ and $\psi(2S)$ at BES

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Many decay modes of $J/\psi$, $\psi(2S)$ and $\chi_{cJ}$ channels are studied using 58M $J/\psi$ and 14M $\psi(2S)$ data sample collected at BES-II. For $J/\psi$ results, two new structures were observed and the invisible decay of $\eta$ and $\eta'$ via $J/\psi \rightarrow \phi\eta, \phi\eta'$ are searched. For $\psi(2S)$ results, the radiative, hadronic baryon pairs decay and leptonic decay modes are measured and the 12% rule are tested. In addition, the many $\chi_{cJ}$ decay channels are observed and most of them are the first measurements.

1 Introduction

The Beijing Spectrometer (BES) is a general purpose solenoidal at the Beijing Electron Positron Collider (BEPC). BEPC operates in the center of mass energy range from 2 to 5 GeV with a luminosity at the $J/\psi$ energy of approximately $5 \times 10^{30} cm^{-2} s^{-1}$. BES detector is described in Ref. 1.

The $J/\psi$ and $\psi(2S)$ has been useful for searches for new hadrons and studies of light hadron spectroscopy. A number of new structures have been observed in $J/\psi$ decays. After observing the strong near-threshold mass enhancements in the $pp$ invariant mass from $J/\psi \rightarrow \gamma pp$ decays 2, the $p\Lambda$ and the $K^+\Lambda$ mass enhancement in $J/\psi \rightarrow pK^+\Lambda$ decays are also observed 3. Recently, we observe the $\omega\phi$ mass enhancement in the double-OZI suppressed decay $J/\psi \rightarrow \gamma\omega\phi$ 4 and $K^+K^-$ enhancement in $J/\psi \rightarrow K^+K^-\pi^0$ 5. The search of invisible decay mode for $\eta$ and $\eta'$ were performed via $J/\psi \rightarrow \phi\eta, \phi\eta'$ 6. The $\psi(2S)$ hadronic decay, $\psi(2S) \rightarrow p\Lambda\pi^0, \eta, p\eta\pi^0, p\eta\pi^0\eta$, and $K^+K^-\pi^0\pi^0$ were measured. The radiative decay, $\psi(2S) \rightarrow \gamma pp, \gamma\pi^+\pi^-\pi^0\pi^0, \gamma\pi^+\pi^-\pi^0\pi^0\eta, \gamma\pi^+\pi^-\pi^0\pi^0\eta$, and $\gamma\pi^+\pi^-\pi^0\pi^0\eta$ were measured with the invariant mass of hadrons below 2.9 GeV have been studied 10. The four baryon decay channels of $\psi(2S) \rightarrow BB$ and one leptonic channel, $\psi(2S) \rightarrow \tau^+\tau^-$ were measured with higher precision 11. The pQCD "12% rule" which is expressed as following:

$$Q_h = \frac{B(\psi(2S) \rightarrow h)}{B(J/\psi \rightarrow h)} = \frac{B(\psi(2S) \rightarrow e^+e^-)}{B(J/\psi \rightarrow e^+e^-)}$$

and leptonic hypothesis are tested. In addition, $\chi_{cJ} \rightarrow K^+K^-K^-\pi^+, \pi^+\pi^-\pi^+\pi^-pp, K\pi\pi^0\pi^0$ and $\eta\pi^+\pi^0\pi^0\eta$ were measured 13,16,17 and many new intermediate states in these channels are observed.

2 Observation of $X(1580)$ and $X(1810)$

Figure 1 shows the Dalitz plot of $M_{h\pi}$ versus $M_{h\pi}$ with the $\gamma\omega\phi$ signal events. Figure 2 shows the distribution of $M_{h\pi}$. The events of upper right in figure 1 corresponding to the bump near
the $\omega \phi$ mass threshold in figure 2. This bump is a new structure, $X(1810)$. PWA gives the mass $1812^{+19}_{-20} \pm 18 MeV/c^2$, width $105 \pm 20 \pm 28 MeV/c^2$ and $J^{PC} = 0^{-+}$, respectively. The combined branching ratio is $B(J/\psi \to \gamma X(1810)) \cdot B(X \to \omega \phi) = (2.61 \pm 0.27 \pm 0.65) \times 10^{-1}$. The possible interpretations for this structure is multiquark, hybrid or glueball.

2.1 $J/\psi \to K^+ K^- \pi^0$

Figure 3 shows Dalitz plot distribution for $M_{K^+ K^- \pi^0}$ versus $M_{K^+ K^- \pi^0}^2$. The events of the upper right in figure 3 corresponds to the bump around 1.6 GeV in figure 1. This bump is a new structure $X(1580)$. PWA gives the pole position, $(1576^{+19}_{-18} \pm 98) - i(109^{+11}_{-12} \pm 126) \text{ MeV/c}^2$, $J^{PC} = 1^{--}$ and the combined branching ratio, $B(J/\psi \to X(1580)^0) \cdot B(X \to K^+ K^-) = (8.5 \pm 0.6) \times 10^{-1}$. Since the width of $X(1580)$ is much broader than the other mesons, it is expected to be a multiquark state.

3 search of $\eta, \eta' \to \text{invisible decay}$

A search of $\eta, \eta' \to \text{invisible}$ decay via $J/\psi \to \phi \eta, \eta'$ were performed. Such a search offer a window for one to observe the physics beyond SM. Two upper limit of branching ratio are given. They are

$$\frac{B(\eta \to \text{invisible})}{B(\eta \to \gamma \gamma)} < 1.65 \times 10^{-3} \text{ at } 90\% C.L., \quad \frac{B(\eta' \to \text{invisible})}{B(\eta \to \gamma \gamma)} < 6.69 \times 10^{-2} \text{ at } 90\% C.L.$$

Here, we give the relative measurement since the many systematic error can cancel.
4 \( \psi(2S) \) hadronic decay

Figure 5, 6 show the distribution of recoiling mass of \( \pi^+\pi^- \) and \( \pi^+\bar{p} \) for \( p\pi \) final state in \( \psi(2S) \) decay after events selection, respectively. Two clear bumps correspond to the signal for \( \psi(2S) \rightarrow p\bar{n}\pi^-\bar{p}n\pi^+ \). Table 1 summarizes the measurement results for \( J/\psi \), \( \psi(2S) \) and \( Q_h \) value. Also, the isospin hypothesis, \( B(p\bar{n}\pi^-) : B(pn\pi^+) : B(pn\pi^-) \approx 1 : 2 : 2 \) is tested.

The multi-body decay mode of \( \psi(2S) \rightarrow \pi^+\pi^-\pi^0 K^+K^- \) are measured and some intermediate states are observed for the first time. Table 2 shows the measured results.

5 \( \psi(2S) \) radiative decay

The radiative decay of \( \psi(2S) \rightarrow \gamma pp, \gamma\pi^+\pi^-\pi^-\pi^-, \gamma K^0\pi^+ K^- + c.c., \gamma\pi^-\pi^- K^+ K^-, \gamma K^+K^-K^-K^- \), \( \gamma\pi^+\pi^-\pi^0 K^+K^- \) and \( \gamma\pi^+\pi^-\pi^0 K^+K^- \) have been studied with the invariant mass of hadrons below 2.9 GeV/c^2. Especially, the \( pp \) mass threshold enhancement structure \( X(1860) \) in \( \psi(2S) \rightarrow \gamma pp \) were searched. Compared with \( J/\psi \rightarrow \gamma pp \), no obvious \( X(1860) \) signal is found. In addition, the radiative decay of \( \psi(2S) \rightarrow \gamma\eta\pi^+\pi^- K^+K^- \) were also observed.

### Table 1: The measurement results for \( \psi(2S) \rightarrow p\bar{n}\pi^-\bar{p}n\pi^+ \), together with the \( J/\psi \) results.

<table>
<thead>
<tr>
<th>Mode</th>
<th>( B(\psi(2S) \rightarrow X) \times 10^{-4} )</th>
<th>( B(J/\psi \rightarrow X) \times 10^{-4} )</th>
<th>( Q_h(%) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( p\bar{n}\pi^- )</td>
<td>2.45 ± 0.11 ± 0.21</td>
<td>2.02 ± 0.17</td>
<td>12.1 ± 1.6</td>
</tr>
<tr>
<td>( \bar{p}n\pi^+ )</td>
<td>2.52 ± 0.12 ± 0.22</td>
<td>1.93 ± 0.13 ± 0.17</td>
<td>13.1 ± 1.8</td>
</tr>
<tr>
<td>( \bar{p}p\pi^0 )</td>
<td>1.32 ± 0.10 ± 0.15</td>
<td>1.09 ± 0.09</td>
<td>12.1 ± 1.9</td>
</tr>
</tbody>
</table>

### Table 2: \( \psi(2S) \rightarrow \pi^+\pi^-\pi^0 K^+K^- \) final state

<table>
<thead>
<tr>
<th>Mode</th>
<th>( B(\psi(2S) \rightarrow X) \times 10^{-4} )</th>
<th>( B(J/\psi \rightarrow X) \times 10^{-4} )</th>
<th>( Q_h(%) )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \pi^+\pi^-\pi^0 K^+ K^- )</td>
<td>11.7 ± 1.0 ± 1.5</td>
<td>120 ± 28</td>
<td>9.8 ± 2.8</td>
</tr>
<tr>
<td>( \omega_0(1710) )</td>
<td>0.59 ± 0.20 ± 0.09</td>
<td>6.6 ± 1.3</td>
<td>8.9 ± 3.8</td>
</tr>
<tr>
<td>( K^*(892)^0 K^-\pi^+\pi^- + c.c. )</td>
<td>8.6 ± 1.3 ± 1.8</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>( K^*(892)^+ K^-\pi^+\pi^- + c.c. )</td>
<td>9.6 ± 2.2 ± 1.7</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>( K^*(892)^0 K^-\rho^- + c.c. )</td>
<td>6.1 ± 1.3 ± 1.2</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>
6 \ \psi(2S) \to BB, \tau^+\tau^-

Four baryon pair, \(p\bar{p}, \Lambda\bar{\Lambda}, \Sigma^0\bar{\Sigma}^0\) and \(\Xi^+\Xi^-\) and a lepton pair are re-measured in \(\psi(2S)\) decay. Our measurement results are very consistent with that of CLEO's\(^{16}\). Besides, the \(\alpha\) value in angular distribution of \(\psi(2S) \to p\bar{p}\) are measured with our large data sample.

According to the leptonic conservation theory, there should be

\[
B(\psi(2S) \to \tau^+\tau^-) \approx B(\psi(2S) \to \mu^+\mu^-) \approx B(\psi(2S) \to \tau^+\tau^-)/0.3885
\]

Table 3 shows our measurement of \(\psi(2S) \to \tau^+\tau^-\), together with the two other leptonic branching ratio. One can find the experimental results is in a good agreement with theoretical prediction for leptonic hypothesis.

7 \ \chi_{cJ} \text{ decay}

Many new \(\chi_{cJ}\) decay mode were observed. They include \(\chi_{cJ} \to \phi K^+K^-, \Sigma^+\Xi^-, \eta\pi^+\pi^-, a_0(980)\pi, \eta f_2(1270)\). Table 4 shows some of measurement results.

10. BES Collaboration, M. Ablikim et al., Hep-ex/0612016
Recent Results on non-$D\bar{D}$ decays of $\psi(3770)$ from BES

Hai Long Ma (For BES Collaboration)

BES Collaboration measured the $R$ values at 3.650, 3.6648 and 3.773 GeV, the $R$ values at 68 energy points in the energy region between 3.650 and 3.872 GeV, the resonance parameters of $\psi(3686)$ and $\psi(3770)$, the branching fractions for $\psi(3770) \to D^0\bar{D}^0$, $D^+D^-$, $D\bar{D}$ and non-$D\bar{D}$, and the observed cross sections for some exclusive light hadron final states at 3.773 and 3.650 GeV. These measurements are made by analyzing the data sets collected with the BESII detector at the BEPC collider.

1 Introduction

There is a long standing puzzle that the observed cross section $\sigma_{DD}^{\text{obs}}$ for $D\bar{D}$ production at the $\psi(3770)$ peak is less than the observed cross section $\sigma_{\psi(3770)}^{\text{obs}}$ for $\psi(3770)$ production\(^1\). Precise or direct measurements of the $R$ values at 3.650, 3.6648 and 3.773 GeV\(^2\), the $R$ values in the energy region between 3.650 and 3.872 GeV\(^3\), the resonance parameters of $\psi(3686)$ and $\psi(3770)$\(^4,5\), the branching fractions for the $\psi(3770)$ decays\(^2,4\) and the observed cross sections for more exclusive light hadron final states produced in $e^+e^-$ annihilation at 3.773 and 3.650 GeV\(^6\) are important to understand the discrepancy. In addition, precise measurements of the $R$ values are also important to test the validity of the pQCD calculation in this energy region and to calculate the effects of vacuum polarization on the parameters of the standard model\(^2,3\). In this paper, we report the results of these measurements.

For convenience, we call the data taken at the c.m. (center-of-mass) energies of 3.650, 3.6648 and 3.773 GeV, the data taken at 49 energy points in the energy region between 3.660 and 3.872 GeV in March 2003, the data taken at 68 energy points in the energy region between 3.650 and 3.872 GeV in December 2003 to be the data A, the data B and the data C, respectively.

2 Measurements of $R$ values

With the data A, we measured the lowest order cross sections $\sigma_{\text{had}}^{\text{obs}}$ and the $R$ values ($R = \sigma_{e^+e^-\text{hadrons}}^{0}/\sigma_{e^+e^-\mu^+\mu^-}^{0}$) for inclusive hadron production at 3.650, 3.6648 and 3.773 GeV. These results\(^2\) are summarized in Table 1, where the first error is the combined statistical and point-to-point systematic error, and the second is the common systematic.

<table>
<thead>
<tr>
<th>$E_{\text{c.m.}}$ (GeV)</th>
<th>3.6500</th>
<th>3.6648</th>
<th>3.7730</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sigma_{\text{had}}^{\text{obs}}$ (nb)</td>
<td>18.98 ± 0.20 ± 0.76</td>
<td>18.30 ± 0.27 ± 0.73</td>
<td>27.68 ± 0.27 ± 1.38</td>
</tr>
<tr>
<td>$R_{\text{had}}$</td>
<td>2.26 ± 0.02 ± 0.09</td>
<td>2.31 ± 0.03 ± 0.09</td>
<td>3.75 ± 0.04 ± 0.19</td>
</tr>
<tr>
<td>$R_{uds}$</td>
<td>2.24 ± 0.02 ± 0.09</td>
<td>2.19 ± 0.03 ± 0.09</td>
<td>-</td>
</tr>
<tr>
<td>$R_{uds+\psi(3770)}$</td>
<td>-</td>
<td>-</td>
<td>3.75 ± 0.04 ± 0.19</td>
</tr>
</tbody>
</table>

With the data C, we also measured the continuum $R_{uds}$ below the $D\bar{D}$ production, the $R_{uds(s)+\psi(3770)}(s)$ and the $R_{\text{had}}(s)$ values in $e^+e^-$ annihilation at all of the 68 energy points\(^3\). They are compared with the other measurements in Fig. 1.

3 Measurements of resonance parameters of $\psi(3686)$ and $\psi(3770)$

A better way to measure the branching fractions for $\psi(3770) \to D\bar{D}$ is to simultaneously analyze the energy-dependent cross sections for the inclusive hadron, $D^0\bar{D}^0$ and $D^+D^-$ production in the energy range covering both the $\psi(3686)$ and $\psi(3770)$. We first accurately measured the
resonance parameters of the two resonances \(^4\) by fitting the observed cross sections at 49 energy points from the data B. In addition, we reported precision measurements of the mass, the total width and the partial leptonic width of the \(\psi(3770)\) by further analyzing the measured \(R\) values at 68 energy points \(^5\) from the data C. The fits to the energy-dependent observed cross sections \(\sigma_{\text{had}}^{\text{obs}}\) and the \(R\) values for inclusive hadron production are shown in Fig. 2. The fitted results are summarized in Table 2.

![Figure 2: The fits to the observed cross sections \(\sigma_{\text{had}}^{\text{obs}}\) and the \(R_{\text{uds}}(\psi(3770))\) values versus \(E_{\text{cm}}\), where the lines show the fits, (a) and (b) are from the data B and (c) is from the data C.](image)

<table>
<thead>
<tr>
<th>Res.</th>
<th>Ref.</th>
<th>(M) (MeV)</th>
<th>(\Gamma_{\text{tot}}) (MeV)</th>
<th>(\Gamma_{\text{lept}}) (eV)</th>
<th>(\Delta M) (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\psi(3770))^B</td>
<td>4</td>
<td>3772.2 ± 0.7 ± 0.3</td>
<td>26.9 ± 2.4 ± 0.3</td>
<td>251 ± 26 ± 11</td>
<td></td>
</tr>
<tr>
<td>(\psi(3686))^B</td>
<td>4</td>
<td>3685.5 ± 0.0 ± 0.3</td>
<td>0.33 ± 0.058 ± 0.002</td>
<td>2330 ± 36 ± 110</td>
<td>86.7 ± 0.7</td>
</tr>
<tr>
<td>(\psi(3770))^C</td>
<td>5</td>
<td>3772.4 ± 0.4 ± 0.3</td>
<td>28.6 ± 1.2 ± 0.2</td>
<td>279 ± 11 ± 13</td>
<td></td>
</tr>
</tbody>
</table>

4 Determinations of \(R_{\text{uds}}\) and \(\sigma_{\psi(3770)}^{\text{obs}}\)

Averaging the measured \(R_{\text{uds}}\) values at 3.650 and 3.6648 GeV from the data A listed in Table 1 by the combined statistical and point-to-point systematic error, we obtain \(R_{\text{uds}} = 2.218 ± 0.019 ± 0.089\), where the first error is the combined statistical and point-to-point systematic error, and the second is the common systematic.

Fitting to the energy-dependent \(\sigma_{\text{had}}^{\text{obs}}\) from the data B yields \(R_{\text{uds}} = 2.262 ± 0.054 ± 0.109\) in the energy region between 3.660 and 3.872 GeV. Fitting to the measured \(R\) values at 68 energy
points from the data C yields \( R_{uds} = 2.121 \pm 0.023 \pm 0.084 \) in the energy region between 3.650 and 3.872 GeV. Here, the errors are statistical and systematic, respectively.

Ignoring the contribution from the continuum \( D\bar{D} \) production at the \( \psi(3770) \) peak, we obtained from the data A to be \( R_{\psi(3770)} = 1.528 \pm 0.042 \pm 0.131 \) due to \( \psi(3770) \) decay into hadrons with the measured \( R_{uds+\psi(3770)} \) value at 3.773 GeV as listed in Table 1 and \( R_{uds} \) measured below \( D\bar{D} \) threshold. This lead to the lowest order cross section \( \sigma^0_{\psi(3770)} \) and the observed cross section \( \sigma^0_{\psi(3770)} \) for \( \psi(3770) \) production at 3.773 GeV. The resonance parameters of the \( \psi(3770) \) obtained by fitting to the energy-dependent \( \sigma^{obs}_{\psi(3770)} \) or fitting to the \( R_{uds+\psi(3770)} \) values from analyzing the cross section scan data samples can further give \( \sigma^0_{\psi(3770)} \) and \( \sigma^0_{\psi(3770)} \) at \( \psi(3770) \) peak. They are summarized in Table 3.

<table>
<thead>
<tr>
<th>Data Sample</th>
<th>Ref.</th>
<th>( \sigma^0_{\psi(3770)} ) [nb]</th>
<th>( \sigma^{obs}_{\psi(3770)} ) [nb]</th>
</tr>
</thead>
<tbody>
<tr>
<td>A</td>
<td>2</td>
<td>9.323 ± 0.103 ± 0.801</td>
<td>7.179 ± 0.195 ± 0.630</td>
</tr>
<tr>
<td>B</td>
<td>4</td>
<td>9.63 ± 0.66 ± 0.35</td>
<td>6.94 ± 0.48 ± 0.28</td>
</tr>
<tr>
<td>C</td>
<td>5</td>
<td>10.06 ± 0.37 ± 0.43</td>
<td>7.25 ± 0.27 ± 0.34</td>
</tr>
</tbody>
</table>

5 Measurements of branching fractions for \( \psi(3770) \rightarrow D^0\bar{D}^0, D^+D^- \) and non-\( D\bar{D} \)

Assuming that there are no other new structures and effects except the \( \psi(3770) \) resonance and the continuum hadron production in the energy region from 3.70 to 3.86 GeV, we can determine the branching fractions for \( \psi(3770) \rightarrow D^0\bar{D}^0, D^+D^- \) and non-\( D\bar{D} \) with the measured \( \sigma^{obs}_{\psi(3770)} \) and the observed cross sections, \( \sigma^{obs}_{D^0\bar{D}^0}, \sigma^{obs}_{D^+D^-} \) and \( \sigma^{obs}_{D\bar{D}} \) for \( D^0\bar{D}^0, D^+D^- \) and \( D\bar{D} \) production measured by analyzing the same data sample A. Fitting to the energy-dependent \( \sigma^{obs}_{D^0\bar{D}^0}, \sigma^{obs}_{D^+D^-} \) and \( \sigma^{obs}_{D\bar{D}} \) in the energy range covering both the \( \psi(3686) \) and \( \psi(3770) \) from the data B, we can also measured the branching fractions for \( \psi(3770) \rightarrow D^0\bar{D}^0, D^+D^- \) and non-\( D\bar{D} \). The measured branching fractions are summarized in Table 4.

![Figure 3: The observed cross sections versus \( E_{cm} \) with fits, where (a) shows inclusive hadron cross section, (b) and (c) show the \( D^0\bar{D}^0 \) and \( D^+D^- \) cross sections, respectively.](image)

<p>| ( \psi(3770) \rightarrow ) | ( \psi(3770) \rightarrow ) | ( \psi(3770) \rightarrow ) | ( \psi(3770) \rightarrow ) |</p>
<table>
<thead>
<tr>
<th>( D^0\bar{D}^0 )</th>
<th>( D^+D^- )</th>
<th>( DD )</th>
<th>non-( D\bar{D} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( B^A[%] )</td>
<td>2</td>
<td>4</td>
<td></td>
</tr>
<tr>
<td>49.9 ± 1.3 ± 3.8</td>
<td>35.7 ± 1.1 ± 3.4</td>
<td>85.5 ± 1.7 ± 5.8</td>
<td>14.5 ± 1.7 ± 5.8</td>
</tr>
<tr>
<td>46.7 ± 4.7 ± 2.3</td>
<td>36.9 ± 3.7 ± 2.8</td>
<td>83.6 ± 7.3 ± 4.2</td>
<td>16.4 ± 7.3 ± 4.2</td>
</tr>
</tbody>
</table>
6 Measurements of the observed cross sections for some exclusive light hadron final states produced in $e^+e^-$ annihilation at 3.773 and 3.650 GeV

We measured the observed cross sections for the exclusive light hadron final states of $\phi\pi^0$, $\phi\eta$, $\phi\pi^+\pi^-$, $\phi K^+K^-$, $\phi\bar{p}\bar{p}$, $2(\pi^+\pi^-)\eta$, $2(\pi^+\pi^-)\pi^0$, $K^+K^-\pi^+\pi^-\pi^0$, $2(K^+K^-)\pi^0$, $p\bar{p}\pi^0$, $p\bar{p}\pi^-\pi^+\pi^0$ and $3(\pi^+\pi^-)\pi^0$ produced in $e^+e^-$ annihilation at $\sqrt{s} = 3.773$ and 3.650 GeV. The preliminary results are shown in Table 5, where the upper limits are set at 90% C.L.. We ignore the interference effects between the continuum and resonance amplitudes, since we do not know the details about the two amplitudes. Therefore we can not draw a conclusion that the $\psi(3770)$ does not decay into these final states even if we do not observe significant difference between the observed cross sections for most light hadron final states at the two energy points.

<table>
<thead>
<tr>
<th>Final State</th>
<th>$\sigma_{e^+e^-\rightarrow\phi\pi^0}^{(ap)}(3.773\text{ GeV})[pb]$</th>
<th>$\sigma_{e^+e^-\rightarrow\phi\pi^0}^{(ap)}(3.650\text{ GeV})[pb]$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\phi\pi^0$</td>
<td>$&lt; 3.5$</td>
<td>$&lt; 8.9$</td>
</tr>
<tr>
<td>$\phi\eta$</td>
<td>$&lt; 12.6$</td>
<td>$&lt; 18.0$</td>
</tr>
<tr>
<td>$\phi\pi^+\pi^-$</td>
<td>$&lt; 11.1$</td>
<td>$&lt; 22.9$</td>
</tr>
<tr>
<td>$\phi K^+K^-$</td>
<td>$15.8 \pm 5.1 \pm 1.8$</td>
<td>$17.4 \pm 9.2 \pm 2.0$</td>
</tr>
<tr>
<td>$\phi\bar{p}\bar{p}$</td>
<td>$&lt; 5.8$</td>
<td>$&lt; 9.1$</td>
</tr>
<tr>
<td>$2(\pi^+\pi^-)\eta$</td>
<td>$153.7 \pm 40.1 \pm 18.4$</td>
<td>$86.6 \pm 40.3 \pm 10.4$</td>
</tr>
<tr>
<td>$2(\pi^+\pi^-)\pi^0$</td>
<td>$80.9 \pm 13.9 \pm 10.0$</td>
<td>$124.3 \pm 21.7 \pm 14.9$</td>
</tr>
<tr>
<td>$K^+K^-\pi^+\pi^-\pi^0$</td>
<td>$171.6 \pm 26.0 \pm 20.9$</td>
<td>$222.8 \pm 37.7 \pm 27.2$</td>
</tr>
<tr>
<td>$2(K^+K^-)\pi^0$</td>
<td>$18.1 \pm 7.7 \pm 2.1$</td>
<td>$&lt; 23.0$</td>
</tr>
<tr>
<td>$p\bar{p}\pi^0$</td>
<td>$10.1 \pm 2.2 \pm 1.0$</td>
<td>$9.2 \pm 3.4 \pm 1.0$</td>
</tr>
<tr>
<td>$p\bar{p}\pi^-\pi^0$</td>
<td>$53.1 \pm 9.2 \pm 6.8$</td>
<td>$29.0 \pm 11.1 \pm 3.7$</td>
</tr>
<tr>
<td>$3(\pi^+\pi^-)\pi^0$</td>
<td>$105.8 \pm 34.4 \pm 16.9$</td>
<td>$126.6 \pm 47.1 \pm 19.2$</td>
</tr>
</tbody>
</table>

7 Summary

Using the data sets collected with the BESII detector at the BEPC collider, BES Collaboration measured the $R$ values at 3.650, 3.6648 and 3.773 GeV, the $R$ values in the energy region between 3.650 and 3.872 GeV, the resonance parameters of $\psi(3686)$ and $\psi(3770)$, the branching fractions for $\psi(3770) \rightarrow D^0\bar{D}^0, D^+D^-, D\bar{D}$ and non-$D\bar{D}$, and the observed cross sections for some exclusive light hadron final states at 3.773 and 3.650 GeV.

References

5. BES Collaboration, M. Ablikim et al., hep-ex/0612056.
VI - Search for New Phenomena
Session

Chairperson: M. Krawczyk
MEASUREMENTS OF B RARE DECAYS AT THE TEVATRON

F. SCURI

on behalf of the CDF and DØ collaborations.

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Largo B. Pontecorvo 3, I-56127 Pisa, Italy

A summary of recent results on B rare decays from the CDF and DØ experiments operating in Run II of the Fermilab Tevatron is given; analyzed decay modes are \( B_{ds} \to hh \), \( B_{ds} \to \mu^+\mu^- \), and \( B \to \mu^+\mu^- h \). Data samples are relative to 1 fb\(^{-1}\) or more integrated luminosity of \( p\bar{p} \) collisions at \( \sqrt{s} = 1.96 \) TeV. All reported results are in agreement with Standard Model predictions and consistent with B-Factories analyzes.

1 Introduction

The large production of all kinds of b-hadrons at the Tevatron offers the opportunity to study rare decays also in the \( B_s \) and b-baryon sectors, exploiting a physics program complementary to the B-Factories.

The bottom anti-bottom production cross-section \( \sigma(b\bar{b}) \) at the Tevatron is \( \mathcal{O}(10^5) \) larger than production in \( e^+e^- \) colliders at the \( \Upsilon(4s) \) or \( Z^0 \) energy scale; however, the inelastic cross-section is a factor \( 10^5 \) larger than \( \sigma(b\bar{b}) \) and the branching ratios of rare b-hadron decays are \( \mathcal{O}(10^{-5}) \) or lower; therefore, interesting events must be extracted from a high track multiplicity environment and detectors need to have very good tracking and vertex resolution, wide acceptance and good particle identification, highly selective trigger. Both CDF and DØ detectors, whose detailed description can be found elsewhere\(^1\), match those requirements.

The following decays modes are analyzed here: charmless \( B_{ds} \to h^+h^- \), \( B_{ds} \to \mu^+\mu^- \), and \( B \to \mu^+\mu^- h \). All those rare decay modes are interesting in the search of New Physics contributions, each mode with its own peculiarity: charmless \( B \to hh \) decays are a useful tool for probing CKM infering on the \( \alpha \) and \( \gamma \) angles of the Unitarity Triangle; they are also sensitive to New Physics effects both in the Penguin diagram contributions to the process and via anomalies in the CP Asymmetry \( A_{CP} \).

The Standard Model prediction of the Branching Ratio for flavor-changing neutral current (FCNC) processes like \( B \to \mu^+\mu^- \) is very suppressed: \( \mathcal{B}(B_s \to \mu^+\mu^-)^{SM} = (3.35\pm0.32) \times 10^{-9} \) (M.Blanke et al.\(^2\)); a slightly higher value for the BR upper limit was recently estimated in Constrained Minimal Flavor Violation\(^3\); \( \mathcal{B}(B_s \to \mu^+\mu^-)^{CM\text{MFV}} < 7.42 \times 10^{-9} \) (95% C.L.); the leptonic branching fraction of the \( B_d \) decay is suppressed by a factor \( |V_{td}/V_{ts}|^2 \) in the CKM matrix elements leading to a SM predicted BR of \( \mathcal{O}(10^{-10}) \). The decay amplitude of \( B_{ds} \to \mu^+\mu^- \) can be significantly enhanced in some extensions of the SM, such as, for instance, the type-II two-Higgs-doublet-model (2HDM), where the BR depends only on \( M_{H^0} \) and on \( \tan\beta \), the ratio between the vacuum expectation values of the two neutral Higgs fields; in this case\(^4\), \( \mathcal{B}(B \to \mu^+\mu^-)^{2HDM} \propto (\tan\beta)^4 \). In the minimal super-symmetric standard model (MSSM) the dependence on \( \tan\beta \) is even stronger, \( \mathcal{B}(B \to \mu^+\mu^-)^{MSSM} \propto (\tan\beta)^6 \), leading to an enhancement of orders of magnitude\(^5\) in case of large values of \( \tan\beta \).

The \( b \to s \mu^+\mu^- \) transition in the \( B \to s \)-meson \( \mu^+\mu^- \) modes allows the study of FCNC in more detail through additional observables, such as the dimuon invariant mass \( m(\mu^+\mu^-) \), and the forward-backward asymmetry of the \( s \)-quark in the dimuon system.

A summary of recent results from the Tevatron experiments is given in the following section for all modes; quoted values refer to integrated luminosities varying from about 1 fb\(^{-1}\) to about 2 fb\(^{-1}\), depending on the specific analysis, and covering the Tevatron Run-II data taking period from 2002 to beginning of 2007.
2 Tevatron result summary for B rare decays

Hadronic decays of the B mesons into two charged tracks are studied in an efficient way by the CDF experiment, exploiting the high performances of its peculiar Secondary Vertex Trigger (SVT), designed to select events with vertexes displaced with respect to the primary vertex and whose full description can be found elsewhere\(^6\). Details of the event selection procedure and of the analysis method can be found in a recent CDF publication\(^7\).

Figure 1: a): M(\(\pi\pi\)) invariant mass distributions; b) Probability Ratio for \(B^0 \rightarrow K^+\pi^-\); c) Probability ratio for \(B_s \rightarrow K^-\pi^+\).

The invariant mass \(m(\pi\pi)\) after cut optimization is shown in figure 1a); about 6500 signal events accumulate in the signal region; individual modes overlap in a single peak as shown by the signal composition (filled colored areas) estimated via Monte Carlo and including full detector simulation. The signal composition in real data is measured with a likelihood fit, combining information from kinematics (mass and momentum) and particle identification; details of the fit procedure can be found elsewhere\(^7\).

<table>
<thead>
<tr>
<th>Decay mode</th>
<th>(BR \times 10^6)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(B^+ \rightarrow \pi^+\pi^-)</td>
<td>5.10 ± 0.33(stat.) ± 0.36(syst.)</td>
</tr>
<tr>
<td>(B^0 \rightarrow K^+K^-)</td>
<td>0.39 ± 0.16(stat.) ± 0.12(syst.)</td>
</tr>
<tr>
<td>(B_s \rightarrow K^+K^-)</td>
<td>24.4 ± 1.4(stat.) ± 4.6(syst.)</td>
</tr>
<tr>
<td>(B_s \rightarrow \pi^+K^-)</td>
<td>5.0 ± 0.75(stat.) ± 1.0(syst.)</td>
</tr>
<tr>
<td>(B_s \rightarrow \pi^+\pi^-)</td>
<td>0.53 ± 0.31(stat.) ± 0.40(syst.)</td>
</tr>
</tbody>
</table>

Table 1: CDF results with 1 fb\(^{-1}\) large sample for \(B(B \rightarrow h^+h^-)\).

Measured branching ratios for individual modes are summarized in table 1; values for the \(B_d\) modes are consistent with the B-Factories results, while the \(B_s\) modes are CDF exclusive.

The very good separation power between \(B^0\) and \(\bar{B}^0\) achieved with the likelihood fit is shown in figure 1b) for the raw asymmetry of the \(B^0 \rightarrow K^+\pi^-\) mode; the probability ratio \(P_R = pdf(B^0)/pdf(B^0) + pdf(B^0)\) is based on probability density functions (pdf) using 5 observables \((M(\pi_1,\pi_2), p_1 + p_2,\) the charged pion momentum um-balancing\(^7\), and the particle identification ID\(_1\) and ID\(_2\)). With this method, the first measurement of direct CP asymmetry \((A_{CP})\) in the \(B_s\) system was done, figure 1e). CDF results on \(A_{CP}\) for the \(\pi\) modes are given in the following equations:

\[
A_{CP}(B_d^0 \rightarrow K^+\pi^-) = \frac{N(B_d^0 \rightarrow K^+\pi^-) - N(B_d^0 \rightarrow K^-\pi^+)}{N(B_d^0 \rightarrow K^-\pi^+) + N(B_d^0 \rightarrow K^+\pi^-)} = -0.086 \pm 0.023(stat.) \pm 0.009(syst.)
\]

\[
A_{CP}(B_s^0 \rightarrow K^-\pi^+) = \frac{N(B_s^0 \rightarrow K^-\pi^+) - N(B_s^0 \rightarrow K^+\pi^-)}{N(B_s^0 \rightarrow K^+\pi^-) + N(B_s^0 \rightarrow K^-\pi^+)} = -0.39 \pm 0.15(stat.) \pm 0.08(syst.)
\]

(1)
Tevatron and B-Factories results on $A_{CP}(B_s^0 \rightarrow K^+\pi^-)$ agree; Tevatron confirms the difference in sign w.r.t. $A_{CP}(B^+ \rightarrow K^+\pi^0)$. The measured value of $A_{CP}(B_s^0 \rightarrow K^-\pi^+)$ (1) can be compared with the value obtained with the Lipkin test, based on minimal assumptions, just SM with SU(3) symmetry; in this case, the following equation holds:

$$A_{CP}(B_s^0 \rightarrow K^-\pi^+) = -A_{CP}(B_s^0 \rightarrow K^+\pi^-) \cdot \frac{B(B_s^0 \rightarrow K^+\pi^-)}{B(B_s^0 \rightarrow K^-\pi^+)} \cdot \frac{\tau(B_s^0)}{\tau(B_s^0)} \approx 0.37$$ (2)

In equation 2 the values of $A_{CP}(B_s^0 \rightarrow K^-\pi^+)$ and $B(B_s^0 \rightarrow K^+\pi^-)$ are taken from the Heavy Flavor Averaging Group, $B(B_s^0 \rightarrow K^+\pi^-)$ is the CDF resulted quoted in table 1, and the lifetime ratio is set equal to one.

The two Tevatron experiments adopt similar procedures to search for $B_s \rightarrow \mu^+\mu^-$ events; a blind cut optimization is made using signal Monte Carlo samples and signal sidebands for estimation of the background in the data sample. The $B^+ \rightarrow J/\psi K^+$ mode is used as normalization channel in the BR estimation. Similar discriminating observables (secondary vertex displacement, B pointing angle to the primary vertex, B isolation) are used to construct a Likelihood Ratio ($L_R$); optimal values of $L_R$ for event counting are obtained with different methods. The main difference in the two analyzes is the definition of the signal window in the $\mu - \mu$ invariant mass spectrum (figure 2); a 360 MeV/c$^2$ wide search region including $B_d$ and $B_s$ is defined by D0, two different regions, 120 MeV/c$^2$ wide and centered in the $B_d$ and $B_s$ nominal masses are used by CDF, quoting separate limits; D0 assumes $B(B_d \rightarrow \mu^+\mu^-) \ll B(B_s \rightarrow \mu^+\mu^-)$ and quotes the overall value as a conservative limit on $B(B_s \rightarrow \mu^+\mu^-)$.

![Figure 2: a), b): D0 Likelihood Ratio $L_R$ vs. $\mu - \mu$ invariant mass; blue lines set the $L_R$ optimal value; c) CDF $\mu - \mu$ invariant mass vs. $L_R$; red and blue boxes define the search areas for $B_d$ and $B_s$ respectively.](image)

Tevatron results on $B_s \rightarrow \mu^+\mu^-$ are summarized in table 2; a statistical combination is made by D0 to quote the overall limit from two data taking periods, before (run IIa) and after (run IIb) insertion of a new inner layer in silicon vertex detector.

<table>
<thead>
<tr>
<th>Exp.</th>
<th>Int. Lumin. pb$^{-1}$</th>
<th>$B_s$ events</th>
<th>$B_d$ events</th>
<th>$B(B_s \rightarrow \mu^+\mu^-)$ 90%(95%) C.L.</th>
<th>$B(B_d \rightarrow \mu^+\mu^-)$ 90%(95%) C.L.</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF</td>
<td>780</td>
<td>1.27±0.37</td>
<td>2.45±0.40</td>
<td>$&lt; 8.0 \cdot 10^{-8}(10)$</td>
<td>$&lt; 2.3 \cdot 10^{-8}(3)$</td>
</tr>
<tr>
<td>D0</td>
<td>1300 (run IIa)</td>
<td>0.8 ±0.2</td>
<td>-</td>
<td>$&lt; 7.5 \cdot 10^{-8}(9.3)(*)$</td>
<td>-</td>
</tr>
<tr>
<td></td>
<td>500 (run IIb)</td>
<td>1.5 ±0.5</td>
<td>2</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

CDF and D0 analyzes to reconstruct the $B \rightarrow \mu^+\mu^- h_s$ modes are similar; each signal is normalized to the analogous $B \rightarrow J/\psi h_s$ mode after a blind cut optimization and background estimation from the signal sidebands; details of the procedure can be found elsewhere.
CDF results for the $B_{u,d}$ modes, obtained with 1 fb$^{-1}$, are summarized in table 3; they are in good agreement with those from the B-Factories$^{11}$ and have similar uncertainty.

<table>
<thead>
<tr>
<th>Mode</th>
<th>Evts. in the signal region</th>
<th>Estimated bkg.</th>
<th>$BR \cdot 10^{-6}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$B^+ \rightarrow \mu^+\mu^-K^+$</td>
<td>90</td>
<td>45.3±5.8</td>
<td>0.60±0.15(stat.)±0.04(syst.)</td>
</tr>
<tr>
<td>$B_d \rightarrow \mu^+\mu^-K^{*0}$</td>
<td>35</td>
<td>16.5±3.6</td>
<td>0.82±0.31(stat.)±0.10(syst.)</td>
</tr>
</tbody>
</table>

CDF and D0 limits for $B(B_s \rightarrow \phi \mu^+\mu^-)$ are shown in table 4; CDF uses a Bayesian approach to extract its limit, while D0 confidence interval is constructed within the Feldman-Cousins scheme.

<table>
<thead>
<tr>
<th>Exp. Int.Lum.</th>
<th>Observed events</th>
<th>Expected events</th>
<th>BR</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF 920 pb$^{-1}$</td>
<td>11</td>
<td>3.5±1.5</td>
<td>&lt; 2.4 \cdot 10^{-6} (90% C.L.)</td>
</tr>
<tr>
<td>D0 450 pb$^{-1}$</td>
<td>0</td>
<td>1.6±0.6</td>
<td>&lt; 3.3 \cdot 10^{-6} (90% C.L.)</td>
</tr>
</tbody>
</table>

3 Conclusions

Tevatron is demonstrated to be a good place to study B rare decays offering different possibilities to constrain more and more New Physics in the $B_s \rightarrow \mu^+\mu^-$ and $b \rightarrow s \mu^+\mu^-$ modes; moreover, a physics program complementary to the B-Factories can be exploited in the $B_s \rightarrow h^+h^-$ sector. CDF made the first observation of the $B_s \rightarrow K^-\pi^+$ mode whose direct CP asymmetry appears to be large in agreement with expectation.

Tevatron current values for the limits on $B(B_s \rightarrow \mu^+\mu^-)$ are entering the $10^{-8}$ territory; D0 current best limit is the first Tevatron result with 2 fb$^{-1}$ integrated luminosity. Any signal at Tevatron before its run II program end ($\sqrt{s}Lt \sim 8$ fb$^{-1}$) will be evidence of New Physics.

CDF and D0 are beginning the exploration of the $b \rightarrow s \mu^+\mu^-$ field; upper limits for $B(B_s \rightarrow \phi \mu^+\mu^-)$ obtained with less than 1 fb$^{-1}$ are close to the SM prediction; CDF reported new results on $B(B_{u,d} \rightarrow s \mu^+\mu^-)$ in general agreement with the B-Factories; significantly improved results will come soon.

References

    CDF Collaboration, CDF public note 8543, Nov. 30, 2006
SEARCH FOR SUPERSYMMETRY AT THE TEVATRON

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This paper reviews some of the most recent results from CDF and DØ experiments on searches for supersymmetry (SUSY) at the Tevatron. We focus on searches for chargino/neutralino, stop, sbottom, and long lived massive SUSY particles, on data samples up to $\sim 1$ fb$^{-1}$. No signal was observed, and constraints are set on the SUSY parameter space.

1 Introduction

Supersymmetry is an extension to the standard model (SM) of particle physics that overcomes some of the theoretical problem in the SM by introducing a new symmetry between fermions and bosons. In this model each SM particle has a superpartner with spin differ by 1/2. The symmetry in SUSY is believed to be broken, and several models (mSUGRA, GMSB, AMSB,...) have been developed to describe the breaking mechanism. In some SUSY models the $R -$ parity quantum number is conserved. In this case SUSY particles must be produced in pairs, and the lightest SUSY particle (LSP) is stable. Cosmological constraints require the LSP to be neutral and colorless$^1$. Thus, the LSP interacts only weakly and escapes detection leading to experimental signature of missing transverse energy ($E_T^*$). In the mSUGRA model, the LSP is usually the lightest neutralino ($\tilde{\chi}^0_1$), and in the GMSB model, the LSP is the gravitino ($\tilde{G}$). CDF and DØ search for SUSY based on their unique signatures in the final states. In this paper we report on recent SUSY search results from CDF and DØ based on $\sim 300$ pb$^{-1}$ to $\sim 1$ fb$^{-1}$ of $p\bar{p}$ collision data at $\sqrt{s} = 1.96$ TeV. The limits presented in this paper are all obtained at 95% confidence level (C.L.) .

$^1R -$ parity $= (-1)^{3(B-L)+2s}$
2 Searches for Chargino and Neutralino

The cross section for associated production of chargino ($\tilde{\chi}_1^\pm$) and neutralino ($\tilde{\chi}_2^0$) is expected to be relatively small compared to other SUSY production via strong interactions. However $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ can decay leptonically ($\tilde{\chi}_1^\pm \rightarrow tW^0$, $\tilde{\chi}_2^0 \rightarrow t^\pm t^\pm$) leading to final states with multiple leptons and large $E_T$ from the escaping $t$ and $\tilde{\chi}_1^0$. This is a very clean final state since the SM background contribution is very small. The background is mainly from production of di-boson, $Z/\gamma^* +$jets, $W-$jets, and $t\bar{t}$. CDF and DØ search for $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ production in 3-lepton plus $E_T$ final states, and in two like-sign lepton plus $E_T$ final states, in a data sample of $\sim 1$ fb$^{-1}$. The number of observed data events and expected SM background events, after applying all the selection cuts, are shown in Table 1. The result indicate no evidence of SUSY and both experiments extract the 95% C.L. upper limit production cross sections, which are shown in Figure 1 and 2. The limits are obtained in the mSUGRA framework, with $\tan \beta = 3$, and with no slepton mixing. The CDF's result exclude chargino mass below 130 GeV/c$^2$. The DØ's result exclude chargino mass below 141 GeV/c$^2$, in the scenario where $m(t) \sim m(\tilde{\chi}_2^0)$ such that the leptonic decay is maximally enhanced.

Table 1: The number of observed data events and expected SM background events in the search for associated production of chargino and neutralino. CDF's $ee + l$ and $\mu\mu + l$ channels are performed on samples collected with low $P_T$ threshold. The DØ's like-sign channel only consists of $\mu^+\mu^-$.  

<table>
<thead>
<tr>
<th>Channel</th>
<th>ee + l</th>
<th>$\mu\mu + l$</th>
<th>$e\ell$</th>
<th>$\mu\ell$</th>
<th>$e\mu$</th>
<th>$t^\pm t^\pm$</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF</td>
<td>1</td>
<td>1</td>
<td>1</td>
<td>0.75</td>
<td></td>
<td>1</td>
</tr>
<tr>
<td># Observed</td>
<td>3</td>
<td>1</td>
<td>0</td>
<td>1</td>
<td></td>
<td>13</td>
</tr>
<tr>
<td># SM Expected</td>
<td>0.97 ± 0.28</td>
<td>0.40 ± 0.12</td>
<td>0.75 ± 0.26</td>
<td>1.26 ± 0.27</td>
<td>7.8 ± 1.1</td>
<td></td>
</tr>
<tr>
<td>DØ</td>
<td>1.1</td>
<td>1.1</td>
<td>1.1</td>
<td>1</td>
<td>0.9</td>
<td>1</td>
</tr>
<tr>
<td># Observed</td>
<td>0</td>
<td>2</td>
<td>0</td>
<td>1</td>
<td>0.94$^{+0.10}_{-0.13}$</td>
<td>1.1 ± 0.4</td>
</tr>
<tr>
<td># SM Expected</td>
<td>0.76 ± 0.67</td>
<td>0.32$^{+0.17}_{-0.03}$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Figure 1: The 95% C.L. upper limit on $\sigma \times BR$ for $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ production as function of $\tilde{\chi}_1^\pm$ mass from CDF.

Figure 2: The 95% C.L. upper limit on $\sigma \times BR$ for $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_2^0$ production as function of $\tilde{\chi}_1^\pm$ mass from DØ.

Figure 3: The 95% C.L. excluded region for the squark and gluino pair production search by DØ.
3 Searches for SUSY in Jets and Missing Energy

3.1 Squark and Gluino

At the Tevatron squark ($\tilde{q}$) gluino ($\tilde{g}$) can be pair produced, and their cascade decays can lead to multiple jets and large $E_T$ final state. The main SM backgrounds are from the production of $Z(\rightarrow \nu\nu)+\text{jets}$, $W(\rightarrow l\nu)+\text{jets}$, $tt$, and QCD multijets. The $E_T$ in QCD multijet background is usually a result of jet energy mis-measurement. DØ search for pair production of squark and gluino in the 2, 3, and 4-jet plus $E_T$ final states, which are respectively, optimized for $\tilde{q}\tilde{q}$, $\tilde{g}\tilde{g}$, and $\tilde{g}\tilde{g}$ productions. To select these events in a data sample of 0.96 fb$^{-1}$, cuts on the jets’ transverse energy ($E_T$), the event $H_T = \sum{jets} E_T$, and $E_T$ are applied. No evidence of squark or gluino production is found, and a limit in the mSUGRA framework is set as a function of squark and gluino mass. This is shown in Figure 3. CDF has also performed a similar search in the 3-jet plus $E_T$ final state, using a data sample of 1.1 fb$^{-1}$. Its selection cuts are optimized for different gluino mass ranges. Similarly no evidence of SUSY is observed, and CDF excludes gluino mass up to $m(\tilde{g}) > 380$ GeV/c$^2$ for the case where $m(\tilde{g}) \sim m(\tilde{q})$.

3.2 Stop and Sbottom

In SUSY, one of the stop and sbottom squarks can be much lighter than the other light flavor squarks. This is due to possible large mixing in the left and right handed weak eigenstates. Therefore, stop and sbottom may be produced at a relatively higher rate than the other squarks at the Tevatron. CDF and DØ search for pair production of stop and sbottom in the heavy flavor jets and $E_T$ final state, using data samples of $\sim 300$ pb$^{-1}$. In these searches it is assumed that the stop decay via a flavor changing neutral current loop ($\tilde{t}_1 \rightarrow c \tilde{t}^0_1$), and the sbottom decay in the channel $b_1 \rightarrow b \tilde{X}^0_1$. To select these events, both experiments require two or three jets in the final state, at least one of the jet tagged as a $c$ or $b$ jet, and the event should have large $E_T$ ($E_T > 50$ GeV) and no isolated lepton. After the event selections, the observed data events are consistent with the expected SM backgrounds, which are mainly from $Z(\rightarrow \nu\nu)+$jets, $W(\rightarrow l\nu)+$jets, $tt$, and QCD multijets. The interpretation of the null results from both stop and sbottom searches are presented as a 95% C.L. exclusion region in the mass planes of $m(\tilde{X}^0_1)$ vs $m(\tilde{t}_1)$ (Figure 4), and $m(\tilde{X}^0_1)$ vs $m(\tilde{b}_1)$. Both CDF and DØ exclude stop mass up to $\sim 140$ GeV/c$^2$ at $m(\tilde{X}^0_1) = 55$ GeV/c$^2$. The sbottom mass is excluded up to 222 GeV/c$^2$ by DØ, and 195 GeV/c$^2$ by CDF.

4 Searches for Long Lived SUSY Particles

4.1 Searches for Long Lived Neutrino

CDF has performed a search for massive long lived particles that decay into photons inside the detector. The search is focused on the GMSB model where the neutralino $\tilde{x}_1^0$, which is the next lightest SUSY particle (NLSP), is long lived and decays into a $\gamma$ and a $\tilde{G}$ inside the detector. The $\tilde{G}$ escapes detection and this leads to $E_T$ signature in the final state. The $\gamma$ from the $\tilde{x}_1^0$ decay arrives at the face of the detector at a later time compared to the other photons that are promptly produced at the interaction point. For this search the events are selected with at least one photon with $E_T > 30$ GeV, one or more jet with $E_T > 35$ GeV, and $E_T > 40$ GeV. The arrival time of the photon candidate is measured in the electromagnetic section of the calorimeter, and is corrected for the time-of-flight assuming that the photon is coming from the primary interaction point. The background consists of prompt photons produced in $pp$ collisions, and fake photons from beam halo and cosmic. The signal is searched in the photon time window of $2 - 10$ ns. In this time window two events are observed in the data, and $1.3 \pm 0.7$ events are
expected from the background. Using this result, an exclusion region in the neutralino lifetime vs neutralino mass plane is determined, and it is shown in Figure 5.

### 4.2 Searches for Charge Massive Particles

Some particles predicted in SUSY carry electrical charge and are massive (CHAMPs). These particles are expected to be slow moving and are very penetrating, just like a “slow muon”. They could decay outside the detector is they are long lived. CDF and DØ have searched for signatures of CHAMP particles inside their detectors. In the search the experiments look for events that contain $\mu$-like particles that are slow moving. CDF makes use of its Time-of-Flight (ToF) and tracking detectors to measure the CHAMP candidates’ velocities and momenta, and then determine its masses. DØ identifies the slow moving $\mu$-like particles based on the timing measurements recorded in the muon detector. No evidence of CHAMP is observed in the data of both experiments. Within the SUSY model with one compactified extra dimension $^2$ where the stop is the LSP, CDF excludes the mass of a stable stop up to 250 GeV/$c^2$ (shown in Figure 6). DØ presents its limits in two models. In the GMSB model DØ assumes stau ($\tilde{\tau}$) is the CHAMP particle, and they obtain an upper limit on the production cross section from 0.62 pb to 0.06 pb for various $\tilde{\tau}$ mass. In the stable chargino model (ref) where the chargino $\chi_1^\pm$ is assumed to be the CHAMP particle, DØ excludes the chargino mass up to 174 GeV/$c^2$.

![Figure 4: The 95% C.L. exclusion region in the mass plane of $m(\chi_1^\pm)$ vs $m(\tilde{t}_1)$ for the search of pair production of stop.](image1)

![Figure 5: The 95% C.L. exclusion region in the $\chi_1^0$ lifetime vs $\tilde{\chi}_1^0$ mass plane for long lived $\chi_1^0$ search.](image2)

![Figure 6: The 95% C.L. upper limit stop production cross section as function of the stop mass.](image3)

### 5 Summary

CDF and DØ have searched for SUSY in up to 1 fb$^{-1}$ of data sample, and have yet found evidence of its existence. Some of the limits set on the SUSY model parameters are the world’s best (example the masses of the squarks and gluino). The Tevatron machine and the experiments are now working optimally. We expect more significant improvement in the SUSY searches at the Tevatron with increase in integrated luminosity and with smarter analysis techniques.

### References

NON SUSY SEARCHES AT THE TEVATRON

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Recent searches for non-SUSY exotics in $p\bar{p}$ collisions at a center-of-mass energy of 1.96 TeV at the Tevatron Run II are reported. The emphasis is put on the results of model-driven analyses which were updated to the full Run IIA datasets corresponding to integrated luminosities of about 1 fb$^{-1}$.

1 Introduction

Numerous searches for non Higgs and SUSY extensions of the Standard Model (SM) are conducted at the Tevatron Run II.

In this report we concentrate on the model-driven analyses that were recently updated by the D0 and CDF collaborations to the full Run IIA datasets, representing integrated luminosities slightly in excess of 1 fb$^{-1}$. Details about these analyses can be found in reference 1.

The SM is constructed with the following ingredients: it's a quantum field theory where the matter fields are replicated into three families of quarks and leptons. This field theory is placed into a four-dimensional space-time. Its lagrangian is invariant under the Poincaré group and under the $SU_C(3) \times SU_L(2) \times U_Y(1)$ gauge group. The electroweak symmetry breaking is provided by the Higgs mechanism.

For about three decades all the experimental tests of the SM have shown no significant deviations with respect to its predictions. However the SM leaves many questions unresolved and is clearly not a full and a satisfactory theory. Therefore many ideas have been proposed to try to extend both predictivity and its domain of validity.

Among these ideas is a possible sub-structure of the particles considered as elementary in the SM. This could explain the replication of the quarks and leptons into three families.
Another path is a possible extension of the SM gauge symmetries. This enables to envisage a unification of the three fundamental interactions described by the SM at a very high energy scale, whilst explaining their differences at low energy as results of different symmetry breakings. One can also postulate the existence of extra space dimensions that could explain the hierarchy between the Planck and the electroweak scales as well as the relative weakness of the gravitational interaction with respect to the three other fundamental interactions.

Searches for fermions sub-structure are reported in section 2, section 3 and section 4 contain searches for hints of extended gauge symmetries and of extra space dimensions respectively. All the exclusion limits are given at the 95% confidence level.

2 Fermions Sub-Structure

In this section, we describe searches driven by two types of models that can be related to a fermions sub-structure. Searches for leptoquarks are presented in the first sub-section and the second sub-section contains searches for fermion compositeness.

2.1 Search for Leptoquarks

The leptoquarks (LQ) carry both a lepton and a baryon quantum number. The relevant phenomenological parameters are $M_{LQ}$ mass for the scalar leptoquarks (simply denoted LQ) and in addition two anomalous couplings for the vector leptoquarks (denoted VLQ).

D0 performed a search for second generation scalar $LQ$ in the $LQ_2 + LQ_2 \rightarrow \mu \nu + q' \bar{q}$ channel, in a dataset of $\int L dt = 1.05 fb^{-1}$. Events containing a hard and isolated muon plus jets and large missing transverse energy ($E_T^\text{miss}$) and $H_T$ are selected. The left hand side of figure 1 shows the $LQ_2$ reconstructed mass where the data are background-like. A limit excluding $M_{LQ_2} < 210$ GeV is set in the hypothesis of this semi-leptonic decay of the $LQ_2$ pairs.

CDF analysed 322 $pb^{-1}$ to search for third generation $VLQ$ in the channel $VLQ_3 + VLQ_3 \rightarrow \tau^+ \tau^- b \bar{b}$ with one $\tau$ subsequently decaying into hadrons and the other into leptons. An $H_T$ variable summing up the $E_T^\text{miss}$ and the $p_T$ of each reconstructed object in the studied topology is used to discriminate the signal from the background. Its distribution does not reveal any data excess. This enables to set limits excluding $M_{VLQ_3} < 251$ GeV and $M_{VLQ_3} < 317$ GeV for minimal and Yang-Mills couplings respectively as displayed in the right hand part of figure 1.

![Graph](image1)

**Figure 1:** Distribution of the $M_{LQ_2}$ in a $LQ_2 + LQ_2 \rightarrow \mu \nu + q' \bar{q}$ search performed by D0 (left). Limit on the production cross section of $VLQ_3 + VLQ_3 \rightarrow \tau^+ \tau^- b \bar{b}$ established by CDF (right).

2.2 Search for Leptons or Quarks Compositeness

For the hypothesis of quarks and leptons compositeness as detailed in reference 2, the main parameters are the excited fermion mass $M_f^*$ and the compositeness scale $\Lambda$. 
We report two searches for excited electrons and muons carried by D0 with integrated luminosities of 1 fb$^{-1}$ and 0.38 fb$^{-1}$ respectively. These excited leptons ($\ell'^{\pm}$) are produced by the following contact interaction process: $q\bar{q} \rightarrow \ell'^{\pm}\ell^{\mp}$ and subsequently decay into the $\ell'^{\pm} \rightarrow \gamma + \ell^{\pm}$ mode, leading to $\gamma + \ell^{\pm}\ell^{\mp}$ final states. The analyses essentially consist in selecting events with two hard and isolated leptons plus a hard and isolated photon and to search for a resonance in the $M_{\ell^{\pm}\ell^{\mp}}$ distribution.

Since no excess of data is found with respect to the SM background, exclusion limits are derived in the $A$ and $M_{\ell^+\ell^0}$ plane. For example, for $A = 1$ TeV: $M_{\ell^+\ell^0} > 756$ GeV and $M_{\ell^{+}\ell^{-}} > 618$ GeV.

We also report that in the measurement of the QCD inclusive jet cross section no deviations with respect to the NLO theory prediction is observed, even up to the highest jet $p_T$ ever probed of about 610 GeV. However, no explicit limit on $M_{\ell^+\ell^0}$ is derived from this measurement yet.

3 Extended Gauge Symmetry

If there exists a grand unification of the strong and the electroweak interactions at a high energy scale, then the breakdown of the corresponding gauge group (i.e. $SU(5)$, $SO(10)$, $E_6$,...) down to the SM gauge group occurs through a cascade of symmetry breakings where extra $SU(2)$ and $U(1)$ factors may appear. Such extra gauge group factors predict the existence of new (and heavy since not observed yet) $W$ and $Z$ bosons that we respectively denote $W'$ and $Z'$.  

We present a D0 search for the $W'^{\pm} \rightarrow e^{\pm}\nu$ process in 0.9 fb$^{-1}$ of data. The events are selected if they contain a hard and isolated electron and a significant $E_T$. The transverse mass $M_T(e^{\pm}, E_T)$ distribution displayed on the left hand part of figure 2 is scrutinized especially above 150 GeV. No data excess is found on top of the SM background tail. Therefore a $W'^{\pm}$ with a mass below 965 GeV is excluded.

We present a CDF search for the $Z' \rightarrow e^+e^-$ process using a data sample of 1.29 fb$^{-1}$. The analysis selects events with two hard and isolated electrons with at least one in the central part of the calorimeter and with a matching track. Here the signal region is defined as the tail (above 150 GeV) of the di-electron invariant mass shown on the left hand side of figure 2. The data are in good agreement with the SM background causing a $Z'$ with a mass lower than 923 GeV and SM-like couplings to be excluded.

![Figure 2: Distributions of the transverse mass (left) and of the invariant mass (right) in searches for $W'$ by D0 and for $Z'$ by CDF respectively.](image)

4 Extended Number of Space Dimensions

The lack of precision measurements of gravity in the sub-millimeter domain leaves some room for possible departures from the Newton's law in this distance range. Such exotic behaviours are obviously predicted in the assumption that space has more than three dimensions because of the
Gauss law. This explains the apparent weakness of the gravitational interaction with respect to the other fundamental interactions.

The Randall-Sundrum (RS) model\(^3\) postulates the existence of a fifth dimension separating two branes. The SM fields are localized on one brane. Gravity lives on the second brane where it isn’t weak, but can propagate along the fifth dimension that has a warped metric. Kaluza-Klein excitations of the gravitons appear as narrow resonances. The relevant parameters are the mass of the first graviton excitation and the $\kappa/\sqrt{M_{Pl}}$ coupling.

The Arkani-Hamed, Dimopoulos and Dvali (ADD) model\(^4\) also localizes the SM fields on one brane and allows gravity to propagate within a bulk possibly made of up to $N = 7$ large extra dimensions. Here Kaluza-Klein excitations of the gravitons cannot be resolved. And the parameters are the number of extra dimensions $N$ and the effective Planck scale $M_D$ (i.e. the Planck scale in 4+$N$ dimensions).

CDF recycled its $Z' \rightarrow e^+e^-$ search into searches for $G \rightarrow e^+e^-$ and $G \rightarrow \gamma\gamma$. The combination of the two latest yields the exclusion plane displayed at the left hand side of figure 3.

We also report on a CDF search in a 1.1$fb^{-1}$ dataset for the $q\bar{q} \rightarrow G + g$ contact interaction process. This process leads to a monojet topology. Hence the analysis requires events with a very hard jet contained in the central part of the calorimeter and confirmed by tracks. In order to allow for a gluon ISR or FSR a second soft jet is accepted. The discriminating variable is the $H_T$ which is compatible with the expected background. Consequently exclusion limits are derived as a function of the number of extra dimensions and the effective Planck scale as shown at the right hand side of figure 3.

![Figure 3: Exclusion plot from the RS CDF search combining the $G \rightarrow e^+e^-$ and $G \rightarrow \gamma\gamma$ channels (left). Limits on large extra dimensions from the CDF monojet search (right).](image)

5 Conclusion

Many searches covering very different topologies have been studied at the Tevatron Run II by the D0 and CDF collaborations. Despite the recent updates of some of these analyses to the full Run IIA datasets of about 1$fb^{-1}$, no hints of exotic extensions of the SM have been found and more stringent exclusion limits have been derived.

References

1. [http://www-d0.fnal.gov/Run2Physics/WWW/results/np.html](http://www-d0.fnal.gov/Run2Physics/WWW/results/np.html)
ELECTROWEAK SYMMETRY BREAKING WITHOUT A HIGGS BOSON
AT THE LHC

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We present two studies into strong symmetry breaking scenarios at the LHC. The first case is
a study into vector boson scattering at ATLAS. This uses the framework of the Electroweak
Chiral Lagrangian with Padé unitarisation to generate possible signal scenarios. Signals could
be observed with an integrated luminosity of \( \int L dt \approx 30 \text{ fb}^{-1} \). Secondly, a search for the tech-
nirho, \( \rho_{TC} \), at CMS is presented, within the Technicolour “Straw Man” model. 5\( \sigma \) discovery
is possible starting from \( \int L dt \approx 4 \text{ fb}^{-1} \).

1 Introduction

It is possible that the higgs boson does not exist, and that a weakly-coupled model is not responsibly
electroweak symmetry breaking. An alternative is that electroweak symmetry breaking
results from new strong interactions. Since the Goldstone bosons resulting from spontaneous
symmetry breaking become the longitudinal components of the \( W \) and \( Z \) bosons at high energy,
we can probe the electroweak symmetry breaking sector by studying vector boson interactions.

Strong electroweak symmetry breaking scenarios can be treated quite generally by an effective
Lagrangian approach, using the Electroweak Chiral Lagrangian accompanied by some
unitarity constraints. A study of vector boson scattering using this framework at ATLAS is
presented in section ???. Under the more specific Technicolour “Straw Man” model, a search for
the technirho, \( \rho_{TC} \), at CMS is presented in section ???.

2 Electroweak Chiral Lagrangian Studies at ATLAS

The Electroweak Chiral Lagrangian \( ^{7} \) (EWChL) describes electroweak interactions at energies
less than 1 TeV. It is built as an expansion in the Goldstone boson momenta. If it is assumed
that custodial symmetry is conserved, there are only two, dimension-4, terms that describe the
quartic couplings of the longitudinal vector bosons

\[ \mathcal{L}^{(4)} = a_4 \left( \text{Tr} \left( D_\mu U D^\mu U^\dagger \right) \right)^2 + a_5 \left( \text{Tr} \left( D_\mu U D^{\mu \dagger} U^\dagger \right) \right)^2 \]

where the Goldstone bosons \( \omega_a \) (\( a = 1, 2, 3 \)) appear in the group element \( U = e^{i \omega_a \sigma_a / v} \), \( \sigma \) are the
Pauli matrices and \( v = 246 \text{ GeV} \). Hence the low-energy effect of the underlying physics in vector
boson scattering is parameterised by the coefficients \( a_4 \) and \( a_5 \).

The Lagrangian does not respect unitarity. To extend its validity range to the higher energies
that we will be probing at the LHC, a unitarisation procedure must be imposed, which can lead to
resonances developing in $[\alpha_s, \alpha_s]$ space. This is dependent on the chosen unitarisation procedure; in the work presented here the Padé or Inverse Amplitude method was used\(^7\).

There have been several studies of EWChL signals in vector boson scattering at ATLAS. All seek to exploit the distinct characteristics of the vector boson fusion process. The boson-boson centre-of-mass energy of interest is $\sim 1$ TeV, so the bosons have high-$p_T$. There are two high energy forward tag jets originating from the quarks that emitted the bosons. Since vector bosons are colourless, there is no colour connection between the tag quarks and hence no additional QCD radiation in the central region.

2.1 WW Scattering: $q\bar{q}WW \rightarrow q\bar{q}WW$

An analysis of $WW \rightarrow l\nu qq$ using the ATLAS fast simulation, ATLFAST, to simulate the effects of the detector is presented here\(^7\). Five signal points in $[\alpha_s, \alpha_s]$ space are chosen; after unitarisation these result in a scalar resonance with a mass of 1 TeV (A), a vector resonance of 1.4 TeV (B), a vector of 1.8 TeV (C), a double resonance of a scalar and a vector (D), and a continuum scenario (E). This final no-resonance scenario is the most pessimistic, with a cross-section x branching ratio of 13 fb. Pythia\(^7\), modified to include the EWChL, is used to simulate the signal and the $W+$jets (where $W \rightarrow b\nu$) and $t\bar{t}$ backgrounds.

The leptonically-decaying $W$ is reconstructed from the highest-$p_T$ lepton and the missing transverse energy, $E_T^{miss}$. The lepton 4-momentum, $E_T^{miss}$ and $W$ mass constraint yield a quadratic equation for the $z$-component of neutrino momentum, $p_Z$. The minimum $p_Z$ solution is chosen because it is closest to the true $p_Z$ in the majority of cases. A cut of $p_T > 320$ GeV is made on this $W$ candidate.

The hadronically-decaying $W$ is highly boosted and can be identified as one or two jets. When jets are identified using the $k_T$ algorithm\(^7\), the highest-$p_T$ jet is chosen as the hadronic $W$ candidate. It is required to have $p_T > 320$ GeV and a mass close to $m_W$. A further "subjet" cut is performed. The $k_T$ algorithm is re-run in subjet mode over the constituents of this jet and the scale at which the jet is resolved into two subjets, $y_{12}p_T^2$, is found\(^7\). For a true $W$, this scale is close to $m_W^2$. A cut requiring $1.55 < \log(p_T \sqrt{y_{12}}) < 2.0$ reduces the $W+$jets background.

To reduce the $t\bar{t}$ background, a crude reconstruction of tops is performed by combining either $W$ candidate with any other jet in the event. Events in which the invariant mass of any of these combinations is close to $m_t$ are rejected. The two tag jets are identified as the highest-$p_T$ jets forward and backward of the $W$ candidates, and required to have $E > 300$ GeV and $|\eta| > 2$. The
\( p_T \) of the full system should be zero, so events with \( p_T(WW + \text{tagjets}) > 50 \text{ GeV} \) are rejected. Finally, events containing more than one additional central jet with \( p_T > 20 \text{ GeV} \) are rejected.

The reconstructed \( WW \) mass after all cuts is shown in figure ?? for the five chosen signal scenarios. All signals are observable above the \( W+\text{jets} \) and \( t\bar{t} \) backgrounds with an integrated luminosity of \( \int \mathcal{L} dt \approx 30 \text{ fb}^{-1} \), with the continuum signal achieving a significance of \( s/\sqrt{b} = 4.7 \).

### 2.2 WZ Scattering: \( qqWZ \rightarrow q'q'WZ \)

A 1.2 TeV vector resonance in \( WZ \) scattering with \( WZ \rightarrow jjll \) (which has \( \sigma \times \text{BR} = 2.8 \text{ fb} \)) was investigated using ATLFAST. The analysis considerations are similar to the above \( WW \) study, although a different implementation of cuts is chosen. After all analysis cuts the only significant background is from \( Z+\text{jets} \) production: for \( 100 \text{ fb}^{-1} \), 14 signal events and 3 background events are expected in the peak region. The reconstructed \( WZ \) mass is shown in figure ??.

![Figure 2: Reconstructed WZ mass for WZ → jjll after all cuts for 300 fb⁻¹.](image)

A recent study using the ATLAS full detector simulation verifies this result, and also finds that significant signals can be observed with 100 fb⁻¹ in the \( WZ \rightarrow l\nu qq \) mode and 300 fb⁻¹ in the \( WZ \rightarrow l\nu ll \) mode. Updated \( WW \) and \( WZ \) scattering analyses will be presented in the forthcoming ATLAS "Computing System Commissioning" note to be completed in summer 2007.

### 3 Search for the technirho, \( \rho_{TC} \), at CMS

The original model of Technicolour (TC) is a scaled-up version of QCD; a new set of interactions is introduced with the same physics as QCD, but at an energy scale \( \Lambda_{TC} \approx 200 \text{ GeV} \). The new strong interaction emerging at the electroweak scale is mediated by \( N_{TC}^{2} - 1 \) technigluons. Electroweak symmetry breaking results from the formation of a technifermion condensate, producing Goldstone bosons (the technipions). Three of the technipions become the longitudinal components of the \( W^\pm \) and \( Z \) bosons.

To generate fermion masses, "Extended Technicolour" interactions are introduced, and the technicolour gauge coupling is required to vary more slowly as a function of the renormalisation scale (it is a "walking" rather than a running coupling). The result is that many technifermions are predicted, and the lightest technicolour resonances appear below 1 TeV. Acquiring the correct top quark mass is a further complication; this is achieved by Topcolour-Assisted Technicolour.

The Technicolour "Straw Man" model sets the framework for searching for the lightest bound states, assuming that these can be considered in isolation. Here we present a search for the colour-singlet \( \rho_{TC} \) in this framework using the CMS detector. The analysis considers the channel \( q\bar{q} \rightarrow \rho_{TC} \rightarrow WZ \) for 14 signal points in \( [m(\rho_{TC}),m(\pi_{TC})] \) space. The cleanest decay mode, \( \rho_{TC} \rightarrow WZ \rightarrow l\nu ll \) is chosen. The \( \sigma \times \text{BR} \) for these signals range from 1 fb to 370 fb.

The main backgrounds are from \( WZ \rightarrow l\nu ll \) and \( ZZ \rightarrow llll, Zbb \rightarrow ll+X \) and \( t\bar{t} \). All signals and backgrounds are generated using Pythia. The CMS fast simulation FAMOS is used, with lepton reconstruction efficiencies and resolutions validated against the GEANT-based full detector simulation.

The three highest-\( p_T \) leptons (electrons or muons) in the event are selected. Making appropriate isolation cuts in the initial identification of these lepton candidates is important in
reducing the $Z\ell\ell$ and $t\bar{t}$ backgrounds. The $Z$ is reconstructed from two same flavour opposite sign leptons. The $W$ is reconstructed from the third lepton and $E_T^{miss}$, as explained in section ??.

Kinematic cuts on the $W$ and $Z$ candidates are needed to improve the signal to background ratio. The $W$ and $Z$ candidates are each required to have $p_T > 30$ GeV. A $Z$ mass window cut of $|m_{t\ell} - m_Z| < 3\sigma$ is particularly effective in reducing the $t\bar{t}$ background. Finally, a cut on the pseudorapidity difference between the $W$ and $Z$ of $|\eta(Z) - \eta(W)| < 1.2$ is effective in reducing the $WZ$ background, although this remains the largest background after all cuts as shown in figure ??(a).

The expected signal sensitivity is computed using the sum of the reconstructed $pT_C$ mass spectra for the signal and backgrounds, taking into account the statistical fluctuations for a given integrated luminosity. It is assumed that the probability density function is Gaussian for the signal and exponential for the background. The sensitivity estimator is given by $S_C = \sqrt{2\ln(L_{S+B}/L_B)}$, where $L_{S+B}$, the signal plus background hypothesis, and $L_B$, the null hypothesis. The sensitivity is computed for each signal point and the resulting contour plot in $(m(pT_C), \sigma(pT_C))$ space is shown in figure ??(a). $5\sigma$ sensitivities are obtained for integrated luminosities starting from 3 fb$^{-1}$, before accounting for systematic uncertainties. Including the expected systematic uncertainties due to the detector, $5\sigma$ discovery is possible starting from 4 fb$^{-1}$ of data.

References

SUSY HIGGS BOSONS AT THE LHC

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Recent results on MSSM Higgs physics at the LHC are reviewed. The dependence of the LHC discovery reach in the \( b\bar{b}H/A, H/A \rightarrow \tau^+\tau^- \) channel on the underlying SUSY scenario is analysed. This is done by combining the latest results for the prospective CMS experimental sensitivities for an integrated luminosity of 30 or 60 fb\(^{-1}\) with state-of-the-art theoretical predictions of MSSM Higgs-boson properties. The results are interpreted in terms of the parameters governing the MSSM Higgs sector at lowest order, \( M_A \) and \( \tan\beta \). While the higgsino mass parameter \( \mu \) has a significant impact on the prospective discovery reach (and correspondingly the "LHC wedge" region), it is found that the discovery reach is rather stable with respect to variations of other supersymmetric parameters. Within the discovery region a determination of the masses of the heavy neutral Higgs bosons with an accuracy of 1-4\% seems feasible. It is furthermore shown that Higgs-boson production in central exclusive diffractive channels can provide important information on the properties of heavy MSSM Higgs bosons.

1 Introduction

Signatures of an extended Higgs sector would provide unique evidence for physics beyond the Standard Model (SM). While models with an extended Higgs sector often give rise to a relatively light SM-like Higgs boson over a large part of their parameter space, detecting heavy states of an extended Higgs sector and studying their properties will be of utmost importance for revealing the underlying physics.

2 Dependence of the LHC discovery reach on the SUSY scenario

In Ref.\(^1\) the reach of the CMS experiment with 30 or 60 fb\(^{-1}\) for the heavy neutral MSSM Higgs bosons has been analysed focusing on the channel \( b\bar{b}H/A, H/A \rightarrow \tau^+\tau^- \) with the \( \tau \)'s
subsequently decaying to jets and/or leptons. The experimental analysis, yielding the number of events needed for a 5σ discovery (depending on the mass of the Higgs boson) was performed with full CMS detector simulation and reconstruction for the final states of di-τ-lepton decays. The events for the signal and background processes were generated using PYTHIA. The experimental analysis has been combined with predictions for the Higgs-boson masses, production processes and decay channels obtained with the code FeynHiggs, taking into account all relevant higher-order corrections as well as possible decays of the heavy Higgs bosons into supersymmetric particles. The results have been interpreted in terms of the two parameters tan β, the ratio of the vacuum expectation values of the two Higgs doublets of the MSSM, and M_A, the mass of the CP-odd Higgs boson. The variation of the discovery contours in the M_A–tan β plane indicates the dependence of the “LHC wedge” region, i.e. the region in which only the light CP-even MSSM Higgs boson can be detected at the LHC at the 5σ level, on the details of the supersymmetric theory. See Ref. for previous analyses.

![Figure 1: Variation of the 5σ discovery contours obtained from the channel bbφ, φ → τ+τ− → jets in the m_h^max benchmark scenario for different values of μ (left plot). The right plot shows the result in the case where no decays of the heavy Higgs bosons into supersymmetric particles are taken into account.](image-url)

Fig. 1 shows the variation of the 5σ discovery contours obtained from the channel bbφ, φ → τ+τ− → jets in the m_h^max benchmark scenario for various values of the higgsino mass parameter μ. The parameter μ enters via higher-order corrections affecting in particular the bottom Yukawa coupling as well as via its kinematic effect in Higgs decays into charginos and neutralinos. Both effects can be seen in Fig. 1. While the left plot shows the full result, in the right plot no decays of the Higgs bosons into supersymmetric particles are taken into account, so that the right plot purely displays the effect of higher-order corrections. Comparison of the two plots shows that in the region of large tan β (corresponding to larger values of M_A on the discovery contours) the dominant effect arises from higher-order corrections. For lower values of tan β, on the other hand, the modification of the Higgs branching ratio as a consequence of decays into supersymmetric particles yields the dominant effect on the 5σ discovery contours. The largest shift in the 5σ discovery contours amounts up to about Δ tan β = 10. The discovery contours have been shown to be rather stable with respect to the impact of other supersymmetric contributions.

The prospective accuracy of the mass measurement of the heavy neutral MSSM Higgs bosons in the channel bbH/A,H/A → τ+τ− is analysed in Fig. 2. The statistical accuracy of the mass measurement has been evaluated via \( \frac{\Delta M_{\hat{a}}}{M_{\hat{a}}} = \frac{R_{M_{\hat{a}}}}{\sqrt{N_{S}}} \), where \( R_{M_{\hat{a}}} \) is the ratio of the di-τ mass resolution to the Higgs-boson mass, and \( N_{S} \) is the number of signal events (φ = H, A). Fig. 2 shows that statistical experimental precisions of 1–4% are reachable within the discovery region. These results are not expected to considerably degrade if further uncertainties related to background effects and jet and missing E_T scales are taken into account. As discussed in Ref., a %-level precision of the mass measurements could in favourable regions of the MSSM.
Figure 2: The statistical precision of the Higgs-boson mass measurement achievable from the channel $b\bar{b}\phi^*, \phi \to \tau^+\tau^- \to \text{jets}$ in the $m_{\phi} \text{max}$ benchmark scenario for $\mu = -200\text{ GeV}$ (left) and $\mu = +200\text{ GeV}$ (right) is shown together with the $5\sigma$ discovery contour.

parameter allow to experimentally resolve the signals of the two heavy MSSM Higgs bosons.

3 MSSM Higgs bosons in central exclusive diffractive production

Adding forward proton detectors to the CMS and ATLAS experiments (at 220 m and 420 m distance around them) can complement the standard LHC physics menu in an interesting way. In particular, “central exclusive diffractive” (CED) Higgs-boson production, where the outgoing protons remain intact and there is no hadronic activity between them, profits from an angular momentum selection rule$^7$ that permits a clean determination of the quantum numbers of the observed Higgs resonance which will be dominantly produced in a scalar state. Other important features of the CED Higgs-boson production process are a potentially excellent mass resolution (irrespective of the decay mode of the produced particle), a much better signal-to-background ratio than conventional Higgs search channels at the LHC, and the possibility to simultaneously access the $b\bar{b}$, $WW$ and $\tau\tau$ decay modes of the Higgs boson(s).

Figure 3: 5$\sigma$ discovery contours for the $H \to b\bar{b}$ channel in CED production in the $M_A - \tan\beta$ plane of the MSSM. The prospective discovery reach is shown in the $m_{\phi} \text{max}$ benchmark scenario (with $\mu = +200\text{ GeV}$). The assumed luminosities are the combined luminosities recorded in ATLAS and CMS of 60 fb$^{-1}$, 60 fb$^{-1} \times 2$, 600 fb$^{-1}$ and 600 fb$^{-1} \times 2$. The values of the mass of the heavy $CP$-even Higgs boson, $M_h$, are indicated by contour lines. The dark shaded (blue) region corresponds to the parameter region that is excluded by the LEP Higgs searches.

In Ref. $^8$ a detailed investigation of the prospects for the MSSM Higgs-boson channels $h, H \to$
\(b\bar{b}, \tau^+\tau^-, WW^*\) in CED production has been carried out (for other studies in the MSSM, see Ref. 9 and references therein). In CED the heavy \(CP\)-even MSSM Higgs boson \(H\) can be produced and its decay into \(b\bar{b}\) can be utilised. While in the SM the \(BR(H \rightarrow b\bar{b})\) is strongly suppressed for \(M_H \gtrsim 2M_W\) because of the dominant decay into gauge bosons, in the MSSM \(H \rightarrow b\bar{b}\) remains by far the dominant decay mode also for larger masses as long as no decays into supersymmetric particles (or lighter Higgs bosons) are open. The CED Higgs-boson production in the \(b\bar{b}\) channel is therefore important over a much larger mass range than in the SM. As an example, Fig. 3 shows the 5\(\sigma\) discovery contours for the \(H \rightarrow b\bar{b}\) channel in CED production in the \(M_A - \tan\beta\) plane of the MSSM (using the \(m_h^{\text{max}}\) benchmark scenario) \(6\) for different luminosity scenarios. It is found that the CED Higgs-boson production channel can cover an interesting part of the MSSM parameter space at the 5\(\sigma\) level if the CED channel can be utilised at high instantaneous luminosity (which requires in particular to bring pile-up background under control). Since there are good chances in this parameter region to detect the heavy neutral MSSM Higgs bosons also in the conventional search channels (see above), a lower statistical significance may be sufficient for the CED channel, corresponding to a larger coverage in the \(M_A - \tan\beta\) plane. The CED Higgs-boson production channel will provide in this case important information on the Higgs-boson properties and may even allow a direct measurement of the Higgs-boson width \(8\).

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SEARCHES FOR NEW PHYSICS
WITH LEPTONS IN THE FINAL STATE AT THE LHC

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Recent estimations for an inclusive Supersymmetry and Extra Dimension models reach for proton-proton collisions at center of mass energy = 14 TeV at the Large Hadron Collider with an integrated luminosity ranging from 1 to 10 fb\(^{-1}\) are reported. Events with leptons (electrons or muons) in the final state are preferred by searches for New Physics.

In 2008, an integrated luminosity of 0.5-1 fb\(^{-1}\) is expected to be collected by the ATLAS\(^1\) and the CMS\(^2\) experiments at the LHC\(^3\) allowing to perform searches for new phenomena over the Standard Model (SM) background. All searches are determined by a precise measurement of the SM processes, which require a detailed understanding of detector performance, effective reconstruction algorithms, full control of trigger rates and a good knowledge of uncertainty sources. Leptons, mainly electrons and muons, have better reconstruction efficiency and energy resolution than taus, jets and missing transverse energy (\(E_T\)). They also provide a clean triggering and a higher background reduction. Above all, leptons may indicate signatures of new physics, e.g. decays of massive strongly interacting particles where leptons are accompanied by jets and \(E_T\), or decays of new resonances to di-lepton pairs with a high invariant mass. For these signatures, a scan over a wide range of parameter space for a given theoretical model sets discovery limits. Supersymmetry (SUSY) and Extra Dimensions (ED) searches are reported.

1 Inclusive SUSY reach with leptons

Supersymmetry\(^4\) described by models which provide a realistic SUSY-breaking scheme. One of the general approaches is given by a Minimal Supergravity\(^5\) (mSUGRA) model with a number of free parameters reduced to 5 variables: mSUGRA\([m_0, m_{1/2}, A_0, \tan\beta, sgn\mu]\) respectively, a common scalar and fermion mass, a trilinear coupling and Higgs sector parameters at the Grand Unification (GUT) scale. In mSUGRA, assuming R-parity\(^4\), new supersymmetric particles are produced in pairs and the lightest one (LSP) is stable and neutral. At the LHC, the SUSY production is dominated by strongly interacting squarks and gluinos \((m_{\text{SUSY}} \sim m_{\tilde{g}})\), which have long decay cascades with jet emission. The cascade ends with the LSP, which is not detected. Therefore, a generic supersymmetry signature is a multi-jets final state with large \(E_T\). The main backgrounds are QCD and \(t\bar{t}, W, Z\) with QCD-jet associated production processes, estimated by using an exact LO evaluation of partonic matrix elements matched with parton showers at the hard process scale\(^6\). In SUSY cascades, leptons are produced in decays of charginos or neutralinos \((e.g., \tilde{\chi}_\pm^0 \rightarrow \chi_\pm^0 l^\pm, \tilde{\chi}_\pm^0 \rightarrow \tilde{\chi}_0^0 l^\pm)\) and the final state consisting of \(n=1-4\) leptons (+jets+\(E_T\)) can be considered. Pairs of leptons can have the same or opposite sign (SS or OS). An important
A. ATLAS preliminary

Figure 1: The mSUGRA discovery reach in $m_0$-$m_{1/2}$ plane for fixed $A_0=0$, $\tan\beta=10$, $\mu>0$ for 1 fb$^{-1}$ with systematic uncertainties denoted by dash lines for the ATLAS (A) and with uncertainties already included for the CMS (B).

The advantage of using signatures with $n(>0)$-leptons is the suppression of the QCD background. The strategy of searches is following. Different experimental signatures are related to the test points of mSUGRA parameter space (in $m_0$-$m_{1/2}$ plane for fixed $A_0$, $\tan\beta$, $\mu$). Results obtained with the full detector simulation and reconstruction (S&R) are extended using fast S&R to other points of the parameter space. In order to obtain the best signal to background ($S/B$) ratio the SUSY selection cuts are optimised for each point. The discovery reach is set for parameters for which the five standard deviation ($5\sigma$) signal significance is expected.

The ATLAS collaboration updated the mSUGRA scanning analysis for $n=0$-2 lepton channels. The cut optimisation was performed with fast S&R (for $m_0=100$-2000 GeV, $m_{1/2}=100$-1500 GeV, $\tan\beta=5$, $10$, $30$, $50$, $A_0=0$, $\mu>0$) using SUSY sensitive observables: $E_T$, $p_T^{\text{jet}}$, $p_T^{\text{miss}}$, Sphericity. Obtained values replaced adequate requirements of the ATLAS standard SUSY selection: $E_T > \text{max}(100 \text{ GeV}, 0.2 \times M_{eff})$, $M_{eff} = E_T + \Sigma^{i=1}_{4} p_T^{\text{jet}_i}$; $n_{\text{jets}}(p_T^{\text{jet}}>50 \text{ GeV}) \geq 4$; $p_T^{\text{miss}}>100 \text{ GeV}$; Sphericity $>0.2$ and for $n$-lepton modes $p_T^{\text{jet}}>20 \text{ GeV}$; Mass $>100 \text{ GeV}$. Background consisted from the following SM processes: $t\bar{t}$, $W(\rightarrow l\nu x)$, $Z(\rightarrow l\ell', \nu \nu)$, $W(\rightarrow l\nu)+N(0-3)$ jets, $Z(\rightarrow l\ell', \nu \nu)+N(2-5)$ jets, $QCD(N(2-6)$ jets and $QQ+N(0-2)$ jets). Large $E_T$ cut (e.g., $>400 \text{ GeV}$) was found to be efficient in the wide $m_{1/2}$ region, reflecting mass splitting between the heaviest and the lightest SUSY particles. The theoretical and experimental uncertainties were considered. The most significant uncertainties, which could enlarge $B$ by a factor of 2, came from the low parton $p_T$ cut and small renormalization scale. The resulting discovery reach, defined by at least 10 signal events and $S/\sqrt{B} > 5$ for 1 fb$^{-1}$, is shown in Fig.1A. Uncertainties lowered the limit curves by about 50 GeV for all channels. Considering this effect, ATLAS reached up to $M_{SUSY} \sim 1.4 \text{ TeV}$ ($m_{1/2} \sim 700 \text{ GeV}$) with $n=0$,1-lepton channels.

In the CMS experiment, several signatures characteristic for mSUGRA were studied (with full S&R and the SUSY selection optimisation) and summarized in the CMS-PTDR. The discovery reach for 1 fb$^{-1}$ of the integrated luminosity is presented in Fig.1B, where curves show limits when systematic uncertainties were taken into account. The most inclusive channels, the jet+$E_T$ and muon+jet+$E_T$ (equivalent to the ATLAS $n=0$,1 lepton modes) give the best results and allow to probe a SUSY existence to the same level of $M_{SUSY} \sim 1.5 \text{ TeV}$ as ATLAS has obtained. Other signatures appearing in the same range of parameters may confirm the discovery. The mSUGRA parameter can be recovered from different quantity measurements, e.g. the reconstruction of SUSY particle masses if the OS-lepton pair mode is observed.

2 Four-lepton signal from Universal Extra Dimensions

The 4-lepton signature is considered in the context of Universal Extra Dimensions (UED) model. The phenomenology of UED is very similar to SUSY, although the origin of UED comes from the sub-millimeter ED model (of the ADD type). In UED, all SM fields are allowed to
propagate along EDs. Therefore, each SM particle has Kaluza-Klein (KK) excitations with the same spin unlike SUSY particles. In the minimal scenario with one ED: mUED[\(R^{-1},AR\)], where \(R^{-1}\) is a size of the ED and \(AR\) an effective cut-off scale, KK excitations have highly degenerated masses even if radiative corrections are introduced. In a long decay chain (from the KK-quark to the lightest and stable KK-photon) soft leptons and jets are emitted. Therefore, the 4-lepton final state (from \(q^{\mu R} \rightarrow Z^{KK R} q, Z^{KK R} \rightarrow l^{KK} l^{\mp}, l^{KK} \rightarrow \gamma^{KK} l^{\pm}\)) is considered the best to be distinguished from the SM background. The CMS study\(^7\) with the full S&K was performed for four mUED points \(AR=20\) and \(R^{-1}=[300, 500, 700, 900]\) GeV analysed in separated channels 4\(\mu\), 4\(e\), 2\(e\)2\(\mu\). Two same-flavour OS isolated lepton pairs are required in the offline selection and the b-tagging and Z-veto rejection are applied (Fig.2A). The discovery potential of the mUED defined as the integrated luminosity needed to measure a signal with a significance of 5\(\sigma\) is shown in Fig.2B. During the first phase of the LHC data taking, the uncertainties due to the limited understanding of the detector performance may limit the sensitivity below 1 fb\(^{-1}\). Although, the mUED can be probed up to \(R^{-1}=600\) GeV with 1 fb\(^{-1}\).

3 Di-lepton resonances from Z-prime and RS Graviton

A high transverse momentum (\(p_T\)) lepton pair is a signature of new massive resonance. Spin-0 gauge bosons \(Z'\) and spin-2 KK excitations of the graviton \(G^*\) with masses in order of 1 TeV are predicted in many EDs and GUT’s theories. At LHC, the resonances are produced directly and they decay promptly to leptons (\(pp \rightarrow X \rightarrow \mu\mu/ee\)). The peak in the invariant mass distribution may be be observed in tails of the SM background processes and the particle mass may be measured directly. The dominant background arises from the lepton pair production (\(pp \rightarrow Z/\gamma \rightarrow l^+l^-\)) in the Drell-Yan process, whereas contributions from the vector boson pair production (\(ZZ, WZ, WW, tt\)) are significantly smaller (a few % of D-Y) and highly suppressed by selection cuts. The K-factors related to the QCD NNLO calculations are used to correct cross-sections in function of the di-lepton mass for the D-Y and the new boson production. Theoretical uncertainties due to a choice of the PDF function and experimental uncertainties are also considered. Moreover, effects of the misalignment of the muon system in the first (<1 fb\(^{-1}\)) and the late (>100 fb\(^{-1}\)) phase of the LHC data taking are included.

In CMS, new reconstruction algorithms are developed to increase the lepton reconstruction efficiency. For electrons, energy of an isolated electromagnetic super-cluster is corrected due to an energy leakage to the hadronic calorimeter and electronics saturation in the electromagnetic calorimeter. The muon track fitting in the tracker and the muon system is optimised to detect and correct effects of energy lost by very high-\(p_T\) muons.

The results of the CMS analyses\(^7\) obtained with the full signal and background S&K are presented in Fig.3. The \(Z'^{\text{SM}}\) boson originated from the so-called Sequential SM or \(Z'^{T}\) from one of GUT’s models\(^{11}\) can be discovered with 1 fb\(^{-1}\) of integrated luminosity (above the Tevatron
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Figure 3: A. CMS $Z'$ reach as a luminosity required for $5\sigma$ discovery of $Z'$ mass. B. The comparison of RS graviton discovery reach for $10\,\text{fb}^{-1}$ in different channels. C. The RS graviton discovery limits in muon channel.

limit of 1 TeV) up to $M_{Z'} \sim 2.6$ TeV (Fig.3A). Graviton excitations $G^*$ are predicted by the Randall-Sundrum\textsuperscript{12} (RS) model. In this model, only one ED is seen by gravity whereas all SM fields live in the 3-dimensional Universe. Couplings and the width of $G^*$ are determined by a parameter $c = \frac{k}{M_{\text{Planck}}}$, where $k \sim M_{\text{Planck}}$ is a mass scale factor. A comparison of the $G^*$ discovery reach in different channels with $10\,\text{fb}^{-1}$ is shown in Fig.3B. The reach for muons is lower due to resolution. The RS graviton decaying to muons can be discovered with $1\,\text{fb}^{-1}$ data up to $M_{G^*} \sim 2.4$ TeV (Fig.3C). After discovery of a new resonant, the spin determination can be obtained\textsuperscript{7}.

Summary

Leptons provide the cleanest signature for exotic searches with the first LHC data. Therefore, inclusive searches with leptons may yield a verification of theoretical predictions for unknown physics beyond the Standard Model. The SUSY discovery potentials obtained by the ATLAS and CMS experiments yield consistent results and show that production of squarks and gluinos with masses of the order of 1.5 TeV can be probed already with an integrated luminosity of $1\,\text{fb}^{-1}$. With $1\,\text{fb}^{-1}$, the searches for EDs can set discovery limits at the level of: $R^{-1} = 600\,\text{GeV}$ for mUED model, $M_{G^*} \sim 2.4$ TeV for the RS graviton and $M_{Z'} \sim 2.6$ TeV for massive $Z'$ bosons.

References

Observability of R-Hadrons at the LHC

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Abstract

Heavy long lived charged particles are predicted in several theories extending the Standard Model. If such particles are coloured they show up as hadrons. These hadrons can be detected in LHC experiments with the very first data exploiting their unique signatures of slow, high momentum particles.

1 Introduction

Several theoretical models predict the possible existence of new long-lived charged particles behaving in collider experiments as stable particles. The predicted particles can be charged under $U(1)$ gauge group (electrically charged) and/or under $SU(3)$ color group. In the latter case hadronized states are expected to appear and interact in the LHC detectors.

The common feature of this class of models is a signature consisting in an high mass stable particle, typically crossing the whole detector as a muon.

Two theoretical models have been studied in details in LHC experiments as benchmarks for a larger classes of scenarios that predict similar signatures: the Gauge Mediated Supersymmetry Breaking models (GMSB) that predict a stable slepton ($\tau$) and the Split Supersymmetry models (Split SUSY) that predict a long lived gluino. In both cases the new quasi-stable particles are expected to have lifetime high enough to cross the detector and a mass higher than $\sim 100$ GeV.

Here the focus is on the latter class of models, in that case gluinos are the long lived particle and they show up in hadronized states called “R-Hadrons”.

The main feature of Split Supersymmetry models is that the masses of scalars are set to a high energy scale ($\Lambda \gtrsim 10^6$ GeV). The gluinos are hence long lived because, being the Next to Lightest Supersymmetric Particle (NLSP) they can only decay into a neutralino (LSP) via a virtual sfermion.

Other models predicting coloured long-lived particles are:
• SuSy breaking with a boundary condition on an extra dimension
• Stable Kaluza-Klein excitation
• Fourth generation fermions

Different R-hadrons types can be formed from a gluino or a squark. The charge of the resulting R-hadron is determined by the quarks bounded to it. While crossing the matter an R-hadron suffers hadronic interactions in which the light quarks bounded to the heavy parton may change, in this case the electric charge of the R-hadron changes.

In order to understand LHC observability the R-hadrons a first crucial step is to understand the effect of the hadronic interactions that a R-hadron suffer when crossing the LHC detectors. The amount of energy loss and the effects of charge flipping should be studied.

2 Simulation of R-Hadron interactions

Some models have been proposed for R-hadrons matter interaction (a review of different models is available) and some of them have been implemented in the simulation programs used for LEP (Geant3 based) and LHC (Geant4 based) experiments.

Two common starting point are usually assumed in those models. First the heavy parton is considered as completely not interacting, acting as a spectator and a reservoir of kinetic energy. Second, the asymptotic cross section value for high $Q^2$ is tuned on the pion-nucleon cross section, with a rescaling simply based on light quark counting (e.g. 2 quarks in case of gluino-mesons and 3 quarks for gluino-baryons).

The energy loss per interaction as a function of $\gamma$, as computed by the Geant 4 models are shown in figure 1. The typical energy loss per interaction is of the order of 1 GeV at $\gamma = 2$, to be compared with the total kinetic energy $E_{kin} > 100$ GeV (for a particle with mass $m > 100$ GeV). This means that the particle is not going to suffer a hadronic showering when crossing calorimeters, and hence is unlikely to be absorbed in the experiments.

A remarkable effect of the several interactions occurring to a R-hadron crossing the LHC experiments calorimeters is that, according to the heavy-parton type (gluino, stop or anti-stop) they tend to convert to a single hadron type. This is due to the fact that interacting with ordinary matter, made of quarks, final states with more antiquark are less probable. So in case of gluino and stop the R-hadrons will become R-baryons after few interaction, while for antistop the preferred state would be the R-meson.

In figure 1 the fraction of R-hadrons converted as function of the penetration depth in iron is shown. Since the interaction cross section are different in case of R-baryons and R-mesons, a scenario in which stop/antistop are long lived can be distinguished, with respect to the gluino case, by the presence of different energy deposits in calorimeters showing some correlation with the particle charge.

![Figure 1: Energy loss per interaction in two different Geant 4 models (left). Fraction of R-hadrons converted in gluino/stop baryons or anti-stop mesons as a function of the penetration depth in iron (right).](image)
3 LHC production and detection

The LHC production of Split Susy gluinos is like that of ordinary Susy gluinos. The production happens mainly via gluon-gluon interaction, with a high total cross section which goes from $\sigma = 50$ nb for a gluino mass of $m = 100$ GeV to $\sigma = 0.1$ pb at $m = 1$ TeV. This means that in this mass range a large number of gluinos are produced even with the first LHC data ($L \sim 1$ fb$^{-1}$).

The velocity and momentum spectra of gluino for masses of $m = 100, 300, 600$ GeV as generated with Pythia$^{10}$ are shown in figure 2. A good fraction of the gluinos have a velocity in the range $0.3 < \beta < 0.9$. In this range particles are fast enough not to be absorbed in the detector (lower $\beta$ would mean very high ionization energy loss) but also significantly different from ultra-relativistic particles such as ordinary particles at momenta in this range ($P > 30$ GeV).

This difference in $\beta$ can be exploited in order to identify heavy long lived charged particles but also lead to some experimental problems due to the fact that LHC experiments are designed with a very short time response (of order of 25 ns) to separate events of different LHC bunch crossing.

In order to identify such particles we need to measure their momenta and their velocity, so that a mass can be reconstructed and they can be distinguished from Standard Model particles. If the particle is charged the momentum can be easily measured in the LHC tracking detectors. The same detectors may provide information on the amount of ionization energy released ($\dfrac{dE}{dx}$). Inverting the Bethe-Bloch formula in the range $0.1 < \beta < 0.9$ it is possible to reconstruct the particle $\beta$ given $\dfrac{dE}{dx}$. A complementary approach for the measurement of the average $\beta$ of a particle reaching the muon systems of LHC experiment is measuring the time of flight using the Drift Tube detectors used in muon spectrometers of CMS and ATLAS.

It is possible to trigger events containing such particles by using muon triggers that can be fired by the long-lived particle itself. This approach has the advantage of being model independent and is hence the preferred way. Complementary approach may rely on the event to be triggered by other high energy particles in the event producing jets or lepton candidates. The main limit of the muon triggers for slow particles is that a minimum velocity of $\beta \geq 0.6$ is required in order to have the event be assigned to the correct LHC bunch crossing. The bunch cross assignment is then used to read the correct data from other subdetectors. If the event is assigned to the wrong bunch crossing the event may fail the higher level trigger and so it may be lost. Nevertheless an high fraction of events are expected to have gluinos with $\beta > 0.6$ so this trigger is still the more promising one. The charge flipping of R-hadrons may add some additional inefficiency in the trigger process so a full simulation of R-hadrons events including hadronic energy loss and charge flipping has been performed in ATLAS and CMS to study the detectability. The result, obtained with some conservative assumption, is that an overall $\sim 15\%$ efficiency can be obtained in CMS for gluinos with $m = 600$ GeV.

Figure 2: Velocity (left) and transverse momentum (right) for three different gluino masses.
If R-hadrons events can be triggered using muons triggers the main backgrounds are expected to be Standard Model events with muons. The best way to distinguish R-hadrons, as well as any heavy stable charged particle, from muons is measuring its velocity and then reconstructing its mass. The mass measurement itself is also interesting being a free parameter in the various models. It has been proved in CMS that using $\frac{\Delta t}{\Delta x}$ and time of flight it is possible to measure velocity with a few % precision and almost completely separate the signal and backgrounds. The resolution obtained combining the two measurements is shown as a function of the particle velocity in figure 3. The optimal region for mass measurement is for $0.6 < \beta < 0.8$ in which the measurement is less biased and with a better resolution. The mass distribution obtained for a 600 GeV gluino with 0.5 fb$^{-1}$ is also shown in figure 3.

4 Conclusion

Using two different techniques for measurement of $\beta$ it will be possible in LHC experiment to search for heavy stable charged particle, including R-hadrons or similar particle originated by a long lived coloured particle. The two methods for $\beta$ measurement have different backgrounds so it is possible, by combining them, to perform a robust and model independent data analysis. Detailed detector understanding, starting with the detectors commissioning in 2007, is needed to give a precise estimate of the discovery reach of LHC experiments with first data.

For long-lived gluinos the discovery should be possible with the first 1 fb$^{-1}$ up to about $m \sim 1$ TeV; for other models the results should be scaled with the production cross section in a proton-proton interaction at 14 TeV.

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We present a concise review of the recent theoretical progress concerning the standard model calculation of the inclusive radiative $B \to X_s \gamma$ decay. Particular attention is thereby devoted to the calculations of the next-to-next-to-leading order fixed-order $O(\alpha_s^2)$ contributions, non-local $O(\alpha_s m_b)$ power corrections, and logarithmic-enhanced $O(\alpha_s \log m_b)$ cut-effects to the decay rate.

1 Introduction

As a flavor-changing-neutral-current process the inclusive radiative $B$-meson ($\bar{B} = \bar{B}^0$ or $B^-$) decay is Cabibbo-Kobayashi-Maskawa-CKM and loop-suppressed within the standard model (SM) and thus very sensitive to new physics (NP) effects. In order to exploit the full potential of $\bar{B} \to X_s \gamma$ in constraining the parameter space of beyond the SM physics both the measurements and the SM prediction should be known as precisely as possible.

The present experimental world average (WA) which includes the latest measurements by CLEO, Belle, and BaBar$^1$ is performed by the Heavy Flavor Averaging Group$^2$ and reads for a photon energy cut of $E_\gamma > E_{\text{cut}}$ with $E_{\text{cut}} = 1.6$ GeV in the $\bar{B}$-meson rest-frame

$$B(\bar{B} \to X_s \gamma)_{\text{exp}} = (3.55 \pm 0.24^{+0.09}_{-0.10} \pm 0.03) \times 10^{-4}. \quad (1)$$

The total error of the WA is below 8% and consists of $i$ combined statistical and systematic error, $ii$ a systematic uncertainty due to the extrapolation from $E_{\text{cut}} = [1.8, 2.0]$ GeV to the reference value, and $iii$ a systematic error due to the subtraction of the $\bar{B} \to X_s \gamma$ event fraction. At the end of the $B$-factory era the final accuracy of the averaged experimental value is expected to be around 5%.

2 Basic Properties of $\bar{B} \to X_s \gamma$

The $b \to s \gamma$ transition is dominated by perturbative QCD effects which replace the power-like Glashow-Iliopoulos-Maiani (GIM) suppression present in the electroweak (EW) vertex by a logarithmic one. This
mild suppression of the QCD corrected amplitude reduces the sensitivity of the process to high scale physics, but enhances the $\bar{B} \to X_s \gamma$ branching ratio (BR) with respect to the purely EW prediction by a factor of around three. The logarithmic GIM cancellation originates from the non-conservation of the tensor current which is generated at the EW scale by loop diagrams involving $W$-boson and top quarks exchange. The associated large logarithms $L \sim \ln M_W/m_b$ have to be resummed at each order in $\alpha_s$, using techniques of the renormalization group (RG) improved perturbation theory. Factoring out the Fermi constant $G_F$, the $b \to s \gamma$ amplitude receives corrections of $O(\alpha_s^3 L^2)$ at leading order (LO), of $O(\alpha_s^3 L^{-1})$ at next-to-leading order (NLO), and of $O(\alpha_s^3 L^{-2})$ at next-to-next-to-leading order (NNLO) in QCD.

A suitable framework to achieve the necessary resummation is the construction of an effective theory with five active quarks, photons and gluons by integrating out the top quark and the EW bosons. Including terms of dimension up to six in the local operator product expansion (OPE) the relevant effective Lagrangian at a scale $\mu$ reads

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{QCD} \times \text{QED}} + \frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb} \sum_{k=1}^{n} C_k(\mu) Q_k.$$  \hspace{1cm} (2)

Here the first term is the conventional QCD and QED Lagrangian for the light SM particles. In the second term $V_{ij}$ denotes the elements of the CKM matrix and $C_k(\mu)$ are the Wilson coefficients of the corresponding operators $Q_k$ built out of the light fields.

The operators and the numerical values of their Wilson coefficients at $\mu_b \sim m_b$ are given by

\begin{equation}
\begin{aligned}
Q_{1,2} &= (\bar{s} \Gamma c)/(\bar{t} \Gamma b), \\
Q_{3-8} &= \frac{e}{16\pi^2} \sum q \left( \bar{q} \Gamma q \right), \\
Q_7 &= \frac{s_{w}^2}{16\pi^2} \sum q \left( \bar{q} \Gamma q \right) T^a b_{\mu} F_{\mu}^\nu, \\
Q_8 &= \frac{s_{w}^2}{16\pi^2} \sum q \left( \bar{q} \Gamma q \right) T^a b_\mu G_{\mu \nu}^\nu,
\end{aligned}
\end{equation} \hspace{1cm} (3)

where $\Gamma$ and $\Gamma'$, entering both the current-current operators $Q_{1,2}$ and the QCD penguin operators $Q_{3-8}$, stand for various products of Dirac and color matrices.\(^3\) In the dipole operator $Q_7$ ($Q_8$), $e$ ($g$) is the electromagnetic (strong) coupling constant, $q_{\mu, \nu}$ are the chiral quark fields, $F_{\mu \nu}$ ($G_{\mu \nu}^\nu$) is the electromagnetic (gluonic) field strength tensor, and $T^a$ are the color generators.

After including LO QCD effects the dominant contribution to the partonic decay rate stems from charm quark loops that amount to $\sim 158\%$ of the total $b \to s \gamma$ decay amplitude. The top contribution is compared to the one from the charm quark with $\sim 60\%$ less than half and has the opposite sign. Diagrams involving up quarks are suppressed by small CKM factors and lead at the amplitude level to an effect of a mere $\sim 2\%$.

All perturbative calculations of $b \to s \gamma$ involve three steps: $i)$ evaluation of the initial conditions $C_k(\mu)\nu)$ of the Wilson coefficients at the matching scale $\mu_w \sim M_W$ by requiring equality of Green's functions in the full and the effective theory up to leading order in (external momenta)/$M_W$, $ii)$ calculation of the anomalous dimension matrix (ADM) that determines the mixing and RG evolution of $C_k(\mu)$ from $\mu_w$ down to the $B$-meson scale $\mu_b \sim m_b$, and $iii)$ determination of the on-shell matrix elements of the various operators at $\mu_b \sim m_b$. Due to the inclusive character of the $B \to X_s \gamma$ mode and the heaviness of the bottom quark, $m_b \gg \Lambda \sim \Lambda_{\text{QCD}}$, non-perturbative effects arise in the last step only as small corrections to the partonic decay rate.

3 Theoretical Progress in $\bar{B} \to X_s \gamma$

At the NNLO level, the dipole and the four-quark operator matching involves three and two loops, respectively. Renormalization constants up to four loops must be found for $b \to s \gamma$ and $b \to sg$ diagrams with four-quark operator insertions, while three-loop mixing is sufficient in the remaining cases. Two-loop matrix elements of the dipole and three-loop matrix elements of the four-quark operators must be evaluated in the last step.

The necessary two- and three-loop matching was performed in Ref.\(^4\) and Ref.\(^5\). The mixing at three loops was determined in Ref.\(^6\) and at four loops in Ref.\(^7\). The two-loop matrix element of the photonic dipole operator together with the corresponding bremsstrahlung was found in Ref.\(^8\) and subsequently confirmed in Ref.\(^9\). These calculations have been very recently extended to include the full charm quark...
mass dependence. The three-loop matrix elements of the current-current operators were derived in Ref. within the so-called large-\(\beta_0\) approximation. A calculation that goes beyond this approximation employs an interpolation in the charm quark mass. The effect of still unknown NNLO contributions is believed to be smaller than the uncertainty that has been estimated after incorporating the above corrections into the SM calculation. To dispel possible doubts about the correctness of this assumption, calculations of the missing pieces are being pursued.

The most impressive bit of the various NNLO calculations is the one of the four-loop ADM that describes the \(O(a_s^4)\) mixing of the four-quark into the dipole operators. It has involved the computation of more than 20000 four-loop diagrams and required a mere computing time of several months on around 100 CPU’s.

Another crucial part of the NNLO calculation is the interpolation in the charm quark mass performed in Ref. The three-loop \(O(a_s^3)\) matrix elements of the current-current operators contain the charm quark, and the NNLO calculation of these matrix elements is essential to reduce the overall theoretical uncertainty of the SM calculation. In fact, the largest part of the theoretical uncertainty in the NLO analysis of the BR is related to the definition of the mass of the charm quark that enters the \(O(a_s)\) matrix elements \(\langle \gamma | Q_{1,2} | b \rangle\). The latter matrix elements are non-vanishing at two loops only and the scale at which \(m_c\) should be normalized is therefore undetermined at NLO. Since varying \(m_c\), between \(m_c(m_b) \sim 1.25\) GeV and \(m_c(m_b) \sim 0.85\) GeV leads to a shift in the NLO BR of more than 10% this issue is not an academic one.

Finding the complete NNLO correction to \(\langle \gamma | Q_{1,2} | b \rangle\) is a formidable task, since it involves the evaluation of hundreds of three-loop on-shell vertex diagrams that are presently not even known in the case \(m_c = 0\). The approximation made in Ref. is based on the observation that at the physical point \(m_c \sim 0.25 \text{ GeV}\) the large \(m_c \gg m_b\) asymptotic form of the exact \(O(a_s)\) and large-\(\beta_0\) \(O(a_s^2\beta_0)\) result matches the small \(m_c < m_b\) expansion rather well. This feature prompted the analytic calculation of the leading term in the \(m_c \gg m_b\) expansion of the three-loop diagrams, and to use the obtained information to perform an interpolation to smaller values of \(m_c\) assuming the \(O(a_s^2\beta_0)\) part to be a good approximation of the full \(O(a_s^3)\) result for vanishing charm quark mass. The uncertainty related to this procedure has been assessed in Ref. by employing three ansätze with different boundary conditions at \(m_c = 0\). A complete calculation of the \(O(a_s^3)\) corrections to \(\langle \gamma | Q_{1,2} | b \rangle\) in the latter limit or, if possible, for \(m_c \sim 0.25 \text{ GeV}\) would resolve this ambiguity and should therefore be attempted.

Combining the aforementioned results it was possible to obtain the first theoretical estimate of the total BR of \(B \to X_s\gamma\) at NNLO. For the reference value \(E_{cut} = 1.6\) GeV the result of the improved SM evaluation is given by

\[
\mathcal{B}(B \to X_s\gamma)_{SM} = (3.15 \pm 0.23) \times 10^{-4},
\]

where the uncertainties from hadronic power corrections (3%), parametric dependences (3%), higher-order perturbative effects (3%), and the interpolation in the charm quark mass (3%) have been added in quadrature to obtain the total error.

The reduction of the renormalization scale dependences at NNLO is clearly seen in Fig. 1. The most pronounced effect occurs in the case of the charm quark mass renormalization scale \(\mu_c\), that was the main source of uncertainty at NLO. The current uncertainty of 3% due to higher-order effects is estimated from the variation of the NNLO curves. The central value in Eq. (4) corresponds to the choice \(\mu_{\text{ref}} = (160, 2.5, 1.5) \text{ GeV}\). More details on the phenomenological analysis including the list of input parameters can be found in Ref. 12.

It is well known that the OPE for \(B \to X_s\gamma\) has certain limitations which stem from the fact that the photon has a partonic substructure. In particular, the local expansion does not apply to contributions
from operators other than $Q_7$, in which the photon couples to light quarks. While the presence of non-local power corrections was thus foreseen such terms have been studied until recently only in the case of the $(Q_8, Q_8)$ interference. In Ref. the analysis of non-perturbative effects that go beyond the local OPE have been extended to the enhanced non-local terms emerging from $(Q_7, Q_8)$ insertions. The found correction scales like $\mathcal{O}(\alpha_s/\Lambda/m_b)$ and its effect on the BR was estimated using the vacuum insertion approximation to be $\sim 0.3, 3.0\%$. A measurement of the flavor asymmetry between $B^0 \rightarrow X_s \gamma$ and $B^- \rightarrow X_s \gamma$ could help to sustain this numerical estimate. Potentially as or maybe even more important than the latter correction are those arising from the $(Q_{1,2}, Q_7)$ interference. Naive dimensional analysis suggests that some non-perturbative corrections to them also scale like $\mathcal{O}(\alpha_s/\Lambda/m_b)$. Since at the moment there is not even an estimate of those corrections, a non-perturbative uncertainty of 5% has been assigned to the result in Eq. (4). This error is the dominant theoretical uncertainty at present and thought to include all known $\mathcal{O}(\alpha_s/\Lambda/m_b)$ terms. Calculating the precise impact of the enhanced non-local power corrections may remain notoriously difficult given the limited control over non-perturbative effects on the light cone.

A further complication in the calculation of $B \rightarrow X_s \gamma$ arises from the fact that all measurements impose stringent cuts on the photon energy to suppress the background from other $B$-meson decay processes. Restricting $E_\gamma$ to be close to the physical endpoint $E_{\max} = m_B/2$, leads to a breakdown of the local OPE, which can be cured by resummation of an infinite set of leading-twist terms into a non-perturbative shape function. A detailed knowledge of the shape function and other subleading effects is required to extrapolate the measurements to a region where the conventional OPE can be trusted.

The transition from the shape function to the OPE region can be described by a multi-scale OPE (MSOPE). In addition to the hard scale $\mu_h \sim m_h \sim 5$ GeV, this expansion involves a hard-collinear scale $\mu_{hc} \sim \sqrt{m_h \Delta} \sim 2.5$ GeV corresponding to the typical hadronic invariant mass of the final state $X_s$, and a soft scale $\mu_s \sim \Delta \sim 1.5$ GeV related to the width $\Delta/2 = m_h/2 - E_{\text{cut}}$ of the energy window in which the photon spectrum is measured. In the MSOPE framework, the perturbative tail of the spectrum receives calculable corrections at all three scales, and may be subject to large perturbative corrections due to the presence of terms proportional to $a_\gamma(\sqrt{m_h \Delta}) \sim 0.27$ and $a_\Delta(\Delta) \sim 0.36$.

A systematic MSOPE analysis of the $(Q_7, Q_7)$ interference at NNLO has been performed in Ref. Besides the hard matching corrections, it involves the two-loop logarithmic and constant terms of the jet and soft function. The three-loop ADM of the shape function remains unknown and is not included. The MSOPE result can be combined with the fixed-order prediction by computing the fraction of events $1 - T$ that lies in the range $E_{\text{cut}} = [1.0, 1.6]$ GeV. The analysis in Ref. yields

$$1 - T = 0.07^{+0.10}_{-0.09, \text{pert}} \pm 0.02, \text{had} \pm 0.02_{\text{pmax}},$$

where the individual errors are perturbative, hadronic, and parametric. The quoted value is almost twice as large as the NNLO estimate $1 - T = 0.04 \pm 0.01, \text{pert}$ obtained in fixed-order perturbation theory and plagued by a significant additional theoretical error related to low-scale perturbative corrections. These large residual scale uncertainties indicate a slow convergence of the MSOPE series expansion in the tail region of the photon energy spectrum. Given that $\Delta$ is always larger than 1.4 GeV and thus fully in the perturbative regime this feature is unexpected.

Additional theoretical information on the shape of the photon energy spectrum can be obtained from the universality of soft and collinear gluon radiation. Such an approach can be used to predict large logarithms of the form $\ln(E_{\max} - E_{\text{cut}})$. These computations have also achieved NNLO accuracy and incorporate Sudakov and renormalon resummation via dressed gluon exponentiation (DGE). The present NNLO estimate of $1 - T = 0.016 \pm 0.003, \text{pmax}$ indicates a much thinner tail of the photon energy spectrum and a considerably smaller perturbative uncertainty than reported in Ref. The DGE analysis thus supports the view that the integrated photon energy spectrum below $E_{\text{cut}} = 1.6$ GeV is well approximated by a fixed-order perturbative calculation, complemented by local OPE power corrections. To understand how precisely the tail of the photon energy spectrum can be calculated requires nevertheless further theoretical investigations.

4 Conclusions

The inclusion of NNLO QCD corrections has lead to a significant suppression of the renormalization scale dependence of the $B \rightarrow X_s \gamma$ branching ratio that have been the main source of theoretical uncertainty at NLO. The central value of the SM prediction is shifted downward relative to all previously published
VII - Higgs and single top Session

Chairperson: B. Klima
Searches for non-Standard-Model Higgs Bosons at the Tevatron

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Search for non-Standard-Model Higgs bosons is one of the major goals of the ongoing Fermilab Tevatron run. Large data sets accumulated by the CDF and DØ experiments break new grounds in sensitivity. We review recent Tevatron results on searches for Higgs bosons in Minimal Supersymmetric Model in the multi $b$-jet and $\tau\tau$ final states, as well as a search for fermiophobic Higgs in the multiphoton final state.

Search for the elusive Higgs bosons is among the most exciting physics topics offered by the high-luminosity Run II of the Fermilab Tevatron. In the Standard Model (SM), the Higgs boson $h^0$ is responsible for Electroweak Symmetry breaking (EWSB) and gives rise to fermion and vector gauge boson masses. However, the Higgs sector may be more complicated than the SM predicts. Generic Two-Higgs Doublet Models (2HDM) contain two complex Higgs doublets, which coupled to up and down-type quarks, respectively. This allows for eight degrees of freedom corresponding to the two complex doublet fields $H_u$ and $H_d$. The ratio of vacuum expectation values of the $H_u$ and $H_d$ doublets is traditionally denoted as $\tan \beta$, which is one of the major parameters in 2HDM. The value of $\tan \beta \approx 45$ equal to the ratio of the top and bottom quark masses (measured at the $Z$ pole) is often referred to as the "golden" value.

Perhaps the most successful and economical theoretical realization of 2HDM is the Minimal Supersymmetric Model (MSSM). In MSSM EWSB occurs naturally via radiative corrections. As a result, three out of the above eight degrees of freedom become longitudinal polarizations of the $W^\pm$ and $Z$ bosons, thus giving rise to their masses. The remaining five degrees of freedom correspond to four scalar Higgs bosons: $h^0$, $H^0$, and $H^\pm$, and one pseudoscalar boson, $A^0$. The SM-like $h^0$ boson is expected to be light, with the mass less than $\sim 135$ GeV, whereas the other four Higgses can be much heavier.

At the Tevatron, the MSSM Higgs bosons $A$ and $H$ are mainly produced either via gluon fusion or radiated off $b$-quarks (see Fig. 1). Production cross section of the $H$ and $A$ bosons is typically enhanced compared to the SM Higgs by the $\tan^2 \beta$ factor, which may be quite large.
This enhancement results in high sensitivity to these particles at the Tevatron, up to their masses of \( \sim 200 \text{ GeV} \), which covers a significant fraction of the allowed MSSM parameter space.

For large \( \tan \beta \), where the Tevatron offers maximum sensitivity, the pseudoscalar Higgs \( A \) is nearly degenerate in mass with either \( h^0 \) or \( H^0 \). Therefore, the final states of interest receive contribution from both the pseudoscalar and one of the scalar Higgses. In what follows we will refer to \( A \) and \( H \) or \( h \) collectively as \( \phi \) and won’t distinguish between the scalar and pseudoscalar Higgses. The major decay modes of \( \phi \) are \( b\bar{b} \) (\( \sim 90\% \)) and \( \tau\tau \) (\( \sim 10\% \)).

In a class of 2HDM, couplings of the light Higgs to fermions are suppressed. We refer to such case as a “fermiophobic Higgs.” As fermiophobic Higgs decay to \( b\bar{b} \) is suppressed, for a sufficiently light Higgs boson the major decay channel could be \( \gamma\gamma \), which proceeds via a loop diagram. DØ Collaboration has performed a search for such a fermiophobic Higgs \( h_f \) produced in association with a charged Higgs. The charged Higgs is assumed to always decay into \( h_f + W^{(*)} \), while \( h_f \) in turn decays to two photons with 100% probability. The full process is therefore \( pp \to h_f h^\pm \to h_f h_f W^{(*)} \to 4\gamma + X \). Experimentally, it is sufficient to require at least three photons in the final state, which results in a very low background and maximizes signal efficiency. The analysis is based on \( \approx 0.8 \text{ fb}^{-1} \) of data collected with a suite of electromagnetic (EM) triggers. The following selection is used in the analysis: at least three photons with the transverse energies above 30, 20, and 15 GeV, respectively, located in the central calorimeter (\( |\eta| < 1.1 \)) and passing standard quality requirements. This selection corresponds to five events observed in data with the expected background of 3.5 \( \pm 0.6 \) events, dominated by direct triple-photon production (2.73 \( \pm 0.55 \) events) and multijet and direct photon events with jets misidentified as photons (0.72 \( \pm 0.15 \) events). Backgrounds are further reduced by requiring the transverse momentum of the 3\( \gamma \) system to exceed 25 GeV. This requirement reduces background to 1.1\( \pm 0.2 \) events and leaves no candidates in data. In the absence of a signal, the following limit is set on the \( h_f \) production: \( \sigma(h_f h^\pm) < 25.3 \text{ fb} \) at the 95% confidence level (C.L.). This cross section limit can be interpreted as the mass limit on \( h_f \) as a function of the charged Higgs mass and \( \tan \beta \). Since associated \( h_f h^\pm \) production depends on \( \tan \beta \) only weakly, the mass limits are similar for small and large values of \( \tan \beta \). For example, \( m_h < 66 \text{ (80) GeV} \) have been excluded at the 95% C.L. for the charged Higgs mass \( < 100 \text{ GeV} \) and \( \tan \beta = 3 \text{ (30) \text{ GeV}} \). The limits become less restrictive with the charged Higgs mass increase. For example, they drop to 44 (60) GeV for the charged Higgs mass \( < 150 \text{ GeV} \).

Recently, DØ reported a new result of the search for \( \phi \) production in association with one or two \( b \)-quarks, which corresponds to the \( 3b \) or \( 4b \) final states. An improved \( b \)-tagging technique based on several tagging algorithms combined via a neural net has been used, resulting in approximately 50% higher efficiency compared to the performance of the best single tagging algorithm. The analysis is based on \( \sim 0.9 \text{ fb}^{-1} \) of data. The selection required at least three \( b \)-tagged jets with transverse energies above 40, 25, and 15 GeV. Signal is then searched for by looking for an excess of events in the invariant mass spectrum of the two leading jets in a mass window around the assumed \( \phi \) mass. The main background comes from multijet production and is estimated from an orthogonal data set selected in a similar way to the signal sample, except for the requirement of exactly two \( b \)-tagged jets. Good agreement is observed between the predicted background and data in the invariant mass spectrum of the two leading jets. Limits
on Higgs production cross section as a function of its mass have been set and are shown in Fig. 2a. Using leading-order (LO) cross section calculations, they are interpreted as limits in the $m_H$-$\tan \beta$ plane as shown in Fig. 2b. For low $\phi$ mass the sensitivity of the search approaches the golden value for $\tan \beta$. Note that the limits on $\tan \beta$ are not directly comparable to the published DØ results$^2$ as the latter are based on next-to-LO cross section.

![Graphs](image)

Figure 2: 95% C.L. limits on (a) the MSSM Higgs production cross section and (b) $\tan \beta$ as a function of the Higgs mass. Area above the red line is excluded.

While the branching fraction of $\phi \rightarrow \tau \tau$ decay is roughly an order of magnitude less than that in the $b\bar{b}$ channel, the $\tau\tau$ final state is clearer than the multijet one, thus offering competitive sensitivity to the MSSM Higgs. Both the CDF and DØ Collaborations have recently reported on new searches for supersymmetric Higgs in the $\tau\tau$ channel based on $\sim 1$ fb$^{-1}$ data sets. The CDF analysis requires one of the $\tau$'s to decay leptonically in the electron or muon channel, and the other $\tau$ to decay either hadronically or leptonically. In the case of both $\tau$'s decaying leptonically, one is required to decay into the electron, while the other – in the muon channel. The DØ analysis is performed in the muon-hadron channel only, but offers competitive sensitivity to the CDF analysis mainly due to more sophisticated neural-net-based $\tau$-identification, which allows for a higher signal efficiency and lower background.

While the total background predictions for the CDF analysis in all three channels agree well with the observed number of events, a slight excess of events is seen in the visible $\tau\tau$ invariant mass spectrum at the mass $\sim 160$ GeV. This excess is consistently seen in both lepton-hadron channels and is consistent with the production of a MSSM Higgs with the mass $\sim 160$ GeV and $\tan \beta \sim 50$ (see Fig. 3a). However, as the significance of the excess (taking into account that it could have been observed at any mass) is less than two standard deviations, CDF still reports limits on the MSSM Higgs production in this channel, which are hurt by the observed excess (see Fig. 4a illustrating that the expected sensitivity is significantly better than the observed limits).

Given this exciting hint from CDF, it is very interesting to see the analogous DØ analysis results. Unfortunately, DØ actually sees a deficit of events in the region of the CDF excess (see Fig. 3b, where a 160 GeV Higgs signal is indicated with shading). Translated into limits on the MSSM Higgs, the DØ analysis excludes the MSSM Higgs with the mass of 160 GeV for $\tan \beta > 40$, thus largely disfavoring an interpretation of CDF excess with the MSSM Higgs production.

More data being collected in the ongoing Tevatron run will allow both Collaborations to extend the reach for non-SM Higgs searches even further and either discover the Higgs or set stringent limits on the MSSM parameters.
Figure 3: Visible $\tau\tau$ mass spectrum in the leptonic-hadronic channel for (a) the CDF and (b) the DØ analysis.

Figure 4: Exclusion in the $m_{\Phi}$-$\tan\beta$ plane for (a) the CDF and (b) the DØ analysis.

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References

HJJ PRODUCTION: SIGNALS AND CP MEASUREMENTS

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Higgs boson production in association with two tagging jets will be mediated by electroweak vector boson fusion and by gluon fusion. For the gluon fusion process, analysis of the azimuthal angle correlations of the two jets provides for a direct measurement of the CP-nature of the $Htt$ Yukawa coupling which is responsible for the effective $Hgg$ vertex.

1 Introduction

Higgs boson production in association with two jets has emerged as a promising channel for Higgs boson discovery and for the study of Higgs boson properties at the LHC. Interest has concentrated on vector-boson-fusion (VBF), i.e. the weak process $qq \rightarrow qqH$ which is mediated by $t$-channel exchange of a $W$ or $Z$, with the Higgs boson being radiated off this weak boson. The VBF production cross section measures the strength of the $WWH$ and $ZZH$ couplings, which, at tree level, require a vacuum expectation value for the scalar field. Hence the VBF channel is a sensitive probe of the Higgs mechanism as the source of electroweak symmetry breaking.

Another prominent source of $Hjj$ events are second order real emission corrections to the gluon fusion process. Such corrections were first considered in Ref.\textsuperscript{1,2} in the large top mass limit and have subsequently been evaluated for arbitrary quark masses in the loops which induce the effective coupling of the Higgs boson to gluons\textsuperscript{3}. For a SM Higgs boson, the generic $Hjj$ cross section from gluon fusion can somewhat exceed the VBF cross section of a few pb\textsuperscript{3} and, thus, gluon fusion induced $Hjj$ events should also provide useful information on Higgs boson properties.

In this contribution we focus on the CP properties of the Higgs Yukawa coupling to the top quark, which is given by

$$\mathcal{L}_Y = y_t H \bar{t} t + i y_A \bar{t} \gamma_5 t,$$

(1)

where $H$ and $A$ denote scalar and pseudo-scalar Higgs fields. Top quark loops then induce effective couplings of the Higgs boson to gluons which, for Higgs masses well below $m_t$, can be described by the effective Lagrangian\textsuperscript{1,2}

$$\mathcal{L}_{\text{eff}} = \frac{y_t}{y_c^M} \cdot \frac{\alpha_s}{12\pi v} \cdot \bar{H} G_{\mu\nu}^a G^{a\mu\nu} + \frac{y_A}{y_c^M} \cdot \frac{\alpha_s}{16\pi v} \cdot A G_{\mu\nu}^a G^{a\mu\nu} \cdot \bar{t} \gamma_5 t,$$

(2)

where $G_{\mu\nu}$ denotes the gluon field strength. From the effective Lagrangian emerge $Hgg$, $Hggg$ and also $Hggg$ vertices, which correspond to triangle, box and pentagon top quark loops and which contribute to gluon fusion processes such as $qq \rightarrow qqH$, $gg \rightarrow qgH$ or $gg \rightarrow ggH$. One example for the first process and for the corresponding VBF diagram is shown in Fig. 1.
2 Azimuthal angle correlations

Analogous to the corresponding VBF case\cite{4,5}, the distribution of the azimuthal angle between the two jets in gluon fusion induced $Hjj$ events can be used to determine the tensor structure of the effective $Hgg$ vertex, which emerges from Eq.(2) as

$$T^{\mu \nu} = a_2 (q_1 \cdot q_2 g^{\mu \nu} - q_1^{\mu} q_2^{\nu}) + a_3 \epsilon^{\mu \nu \rho \sigma} q_1^{\rho} q_2^{\sigma}, \quad (3)$$

with $a_2 = \frac{g_s^2}{64 \pi^2} \cdot \frac{1}{2}$ and $a_3 = -\frac{g_s^2}{64 \pi^2} \cdot \frac{1}{2}$. We assume SM-size couplings in our analysis below.

In resolving interference effects between the CP-even coupling $a_2$ and the CP-odd coupling $a_3$ it is important to measure the sign of the azimuthal angle between the jets. Naively one might assume that this sign cannot be defined unambiguously in pp collisions because an azimuthal angle switches sign when viewed along the opposite beam direction. However, in doing so, the “toward” and the “away” tagging jets also switch place, i.e. one should take into account the correlation of the tagging jets with the two distinct beam directions. Defining $\Delta \Phi_{jj}$ as the azimuthal angle of the “away” jet minus the azimuthal angle of the “toward” jet, a switch of the two beam directions leaves the sign of $\Delta \Phi_{jj}$ intact.\cite{5} The corresponding distributions, for two jets with

$$p_T > 30 \text{ GeV}, \quad |\eta_j| < 4.5, \quad |\eta_t - \eta_j| > 3.0,$$ \quad (4)

are shown in Fig. 2 for three scenarios of CP-even and CP-odd Higgs couplings. All three cases are well distinguishable. The maxima in the distributions are directly connected to the size of the parameters $a_2$ and $a_3$ in Eq. (3). For

$$a_2 = a \cos \alpha, \quad a_3 = a \sin \alpha,$$ \quad (5)

the positions of the maxima are at $\Delta \Phi_{jj} = \alpha$ and $\Delta \Phi_{jj} = \alpha \pm \pi$.

3 Observability at the LHC

The azimuthal angle correlations of the two leading jets in gluon fusion are fairly independent of the Higgs boson mass and they do not depend on the Higgs decay mode, except via kinematical effects due to cuts on the decay products. In order to observe them, however, background processes have to be suppressed by a sufficient degree. Clearly, this depends on which decay channels are available for the Higgs boson. The most promising case is a SM-like Higgs boson of mass around $m_H \approx 160$ GeV with decay $H \rightarrow WW^* \rightarrow l^+l^-\nu\bar{\nu}$ ($l = e, \mu$). We here give a brief summary of our findings. Details of the parton level simulation are given Ref.\cite{6} Similar to the analogous $H \rightarrow WW$ search in VBF\cite{7} the dominant backgrounds arise from $t\bar{t}$ production in association with 0, 1, or 2 additional jets and from $WWjj$ production at order $\alpha^2\alpha_s^2$ (QCD $WWjj$ production) or at order $\alpha^4$ (EW $WWjj$ production, which includes $H \rightarrow WW$ in VBF).
Figure 2: Normalized distributions of the jet-jet azimuthal angle difference as defined in the text. The curves are for the SM CP-even case (as = 0), a pure CP-odd (as = 0) and a CP-mixed case (as = as ≠ 0).

Table 1: Signal and background cross sections and the expected number of events for \( \mathcal{L}_{\text{int}} = 30 \text{ fb}^{-1} \) at different levels of cuts.

<table>
<thead>
<tr>
<th>process</th>
<th>inclusive cuts</th>
<th>selection cuts</th>
<th>selection cuts &amp; ( \Delta R_{jj} &gt; 3 )</th>
</tr>
</thead>
<tbody>
<tr>
<td>GF ( pp \to H + jj )</td>
<td>115 fb</td>
<td>31.5 fb</td>
<td>945 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>EW ( pp \to W^+W^- + jj )</td>
<td>75 fb</td>
<td>16.5 fb</td>
<td>495 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>( pp \to t\bar{t} )</td>
<td>6830 fb</td>
<td>23.3 fb</td>
<td>699 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>( pp \to t\bar{t} + j )</td>
<td>9520 fb</td>
<td>51.1 fb</td>
<td>1530 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>( pp \to t\bar{t} + jj )</td>
<td>1680 fb</td>
<td>11.2 fb</td>
<td>336 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>QCD ( pp \to W^+W^- + jj )</td>
<td>363 fb</td>
<td>11.4 fb</td>
<td>342 events / 30 fb(^{-1} )</td>
</tr>
<tr>
<td>sum of backgrounds</td>
<td>18500 fb</td>
<td>114 fb</td>
<td>3410 events / 30 fb(^{-1} )</td>
</tr>
</tbody>
</table>

The first column of Table 1 gives the expected LHC cross sections for the fairly inclusive cuts of Eq. 4 (but \( |\eta_1 - \eta_2| > 1 \)) for the tagging jets, defined as the two highest \( p_T \) jets in an event, and lepton cuts given by

\[
\begin{align*}
    p_T > 10 \text{ GeV}, \quad |\eta_l| < 2.5, \quad \Delta R_{jj} = \sqrt{(\eta_1 - \eta_2)^2 + (\Phi_1 - \Phi_2)^2} > 0.7. 
\end{align*}
\]

The large top quark background can be suppressed by a veto on events with a \( b \)-quark tag on any observable jet. A characteristic feature of \( H \rightarrow WW \) decay is the small angular separation and small invariant mass of the \( l^+l^- \) system, which is exploited by the cuts

\[
\begin{align*}
    \Delta R_{ll} < 1.1, \quad m_{ll} < 75 \text{ GeV}. 
\end{align*}
\]

The signal is further enhanced by requiring large lepton transverse momentum, \( p_T > 30 \text{ GeV} \), a transverse mass of the dilepton/missing \( E_T \) system consistent with the Higgs mass, \( m_{WW}^T < m_H + 10 \text{ GeV} \) and not too small compared to the observed dilepton mass, \( m_{ll} < 0.44 \cdot m_{WW}^T \), and a significant amount of missing \( p_T \), \( \not{p}_T > 30 \text{ GeV} \). The resulting cross sections and expected event rates for \( 30 \text{ fb}^{-1} \) are given in the second and third columns of Table 1: with \( 30 \text{ fb}^{-1} \) the
LHC can establish a Higgs signal in gluon fusion with a purely statistical error leading to a significance of $S/\sqrt{B} = 16$.

The resulting event sample of about 950 signal and 3400 background events is large enough and sufficiently pure to analyze the azimuthal angle between the two tagging jets. One finds, however, that the characteristic modulation of the $\Delta \phi_{jj}$ distribution in Fig. 2 is most pronounced for large rapidity separations of the tagging jets.\(^6\) Imposing $\Delta \eta_{jj} > 3$, one obtains the cross sections in the second to last column of Table 1 and azimuthal angle distributions as shown in Fig. 3 for an integrated luminosity of 300 fb$^{-1}$. Already for 30 fb$^{-1}$ of data, however, can one distinguish the SM expectation in the left panel of Fig. 3 from the CP-odd case in the right panel with a purely statistical power of more than 5 sigma. We do not expect detector effects or higher order QCD corrections to substantially degrade these conclusions.

Figure 3: The $\Delta \phi_{jj}$ distribution for a pure CP-even coupling (left) and a pure CP-odd coupling (right) for $L_{int} = 300$ fb$^{-1}$. From top to bottom: GF signal, EW $W^+W^-jj$, $t\bar{t}$, $t\bar{t}jj$, and QCD $W^+W^-jj$ backgrounds.

Acknowledgments

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References

MONTE CARLO STUDIES OF THE JET ACTIVITY IN HIGGS PRODUCTION IN ASSOCIATION WITH 2 JETS AT THE LHC

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Distributions between two jets in Higgs+2-jet events have been analysed in the past at leading order as a means to probe the Higgs couplings. Here we explore the impact of higher-order corrections on the azimuthal angle correlation of the two leading jets and on the rapidity distribution of extra jets. Our study includes matrix-element and shower MC effects, for the two leading sources of Higgs plus two jet events at the CERN LHC, namely vector-boson and gluon fusion. We show that the discriminating features present in the previous leading-order matrix element studies survive.

1 Higgs + 2-jet production

Finding the Higgs boson and studying its properties will be one of the main tasks of the LHC. The two leading production modes of a SM-like Higgs boson are gluon fusion and vector-boson fusion (VBF). The latter is characterised by two forward tagging jets separated by a large rapidity interval, a feature that is very helpful to suppress backgrounds. Of course also gluon-fusion processes give rise to Higgs plus two jet production. For a measurement of Higgs couplings it is important to distinguish these two sources of Hjj events. Fortunately, the distributions of the two tagging jets, as well as their correlations, are markedly different for gluon fusion and VBF. For example, the dijet invariant mass distribution of the two leading jets in gluon fusion is substantially softer than in VBF. This is due to the different shape of the PDF's of the partons initiating the hard scattering; in VBF the scattering occurs mostly through valence quarks, in gluon fusion through gluons.
2 The azimuthal correlation

A second characteristic difference emerges in the azimuthal correlations of the two tagging jets\(^4\). The distribution of the azimuthal angle \(\Delta\phi_{jj}\) between the jets directly reflects the tensor structure of the coupling of the Higgs boson to weak bosons or gluons\(^5,6\). The SM \(HHW\) and \(HZZ\) couplings lead to a fairly flat \(\Delta\phi_{jj}\) distribution. In contrast, the loop induced effective \(Hgg\) coupling, which, in the large top-mass limit, can be written as a CP-even effective Lagrangian, produces a dip in the \(\Delta\phi_{jj}\) distribution at 90 degrees. The same correlation and similar dynamical properties were used in Refs.\(^5,6\) to discriminate between the SM coupling and anomalous (New Physics) \(HHW\) and \(HZZ\) couplings in VBF processes.

The analyses of Refs.\(^4,5,6\) were done at the parton level only. In Fig. 1 we show a sample of contributing diagrams to the Higgs production in gluon fusion. Previous experience with the azimuthal correlation between two jets at large rapidity intervals in dijet production in \(p\bar{p}\) collisions, analysed at the parton level\(^7,8,9,10,11\) and with parton showers and hadronisation\(^12,13\), and measured at the Tevatron\(^14\), leads us to expect that a certain amount of de-correlation between the jets will be induced by showering and hadronisation. This reduces the correlation induced by the dynamical properties of Higgs + 2 jet production at the parton level. An early study did, indeed, find a much weaker correlation between the tagging jets in Higgs + 2 jet production via gluon fusion after showering and hadronisation\(^15\). This analysis did not allow, though, for a direct comparison with the result of Ref.\(^4\), because in Ref.\(^15\) the two tagging jets associated to the Higgs production were generated by the parton shower and not at the matrix element level. Thus, it was not possible to distinguish the decorrelation due to showering and hadronisation from an inherent lack of correlation between the two tagging jets caused by the approximations in the parton-shower generation. In Ref.\(^16\) we addressed the shortcomings of

![Diagram](image)

Figure 1: Samples of Feynman graphs contributing to \(H + 2\) jet production via gluon fusion. From Ref.\(^4\).

either a purely hard matrix element calculation or of a purely parton-shower approach. By using ALPGEN\(^17,18\) to calculate the matrix elements for emission of hard partons and HERWIG\(^13\) to then evolve the parton-level events through the shower and hadronisation phases we considered the azimuthal correlation between the two tagging jets and a veto on the jet activity in the rapidity interval between the tagging jets. In the case of the azimuthal correlation, we found that the dip at \(\Delta\phi_{jj} = \pi/2\), characteristic of a CP-even Higgs boson produced via gluon fusion\(^4,5\), is slightly filled by the parton shower, but not as much as one would find by generating the tagging jets through the parton shower\(^15\). This feature is shown in Fig. 2, where the dot-dashed line gives the parton-level Lowest Order (LO) prediction, while the solid line represents the distribution including also shower evolution effects. As a reference, the dashed histogram shows the \(\Delta\phi_{jj}\) distribution for the VBF production channel, where the effects of parton showering are almost indistinguishable from the parton level calculation. It is worth mentioning that the shower effects on the \(\Delta\phi_{jj}\) distribution are of the same order as the ones given by the QCD next-to-leading corrections, recently computed in Ref.\(^19\).
Figure 2: Normalized distribution of the azimuthal distance between the two tagging jets in Higgs + 2 parton production via gluon fusion with (solid histogram) and without (dot-dashed curve) parton shower and via VBF with parton shower (dashes).

Figure 3: Left: Normalised distribution of the multiplicity of additional jets which fall within the rapidity interval of the tagging jets, generated via Higgs + 2 (solid) and 3 (dashes) final-state-parton production. The left panel is for Higgs production via gluon fusion, the right panel corresponds to VBF.

3 The central jet veto

A veto of any additional jet activity in the central rapidity region is expected to suppress the backgrounds to VBF processes compared to the signal, because the QCD backgrounds are characterised by quark or gluon exchange in the t-channel. The exchanged partons, being coloured, are expected to radiate off more gluons.

In order to quantify the jet activity in the rapidity interval between the tagging jets, we have computed the multiplicity distribution for jets that fall within this rapidity interval. The multiplicity is normalised to the total cross section for Higgs + 2-jet production after jet reconstruction. In Fig. 3 (left) we show the distribution of the resulting additional jet multiplicity for Higgs + 2 (solid) and 3 (dashes) final-state-parton production via gluon fusion. Analogous results for VBF are shown in Fig. 3 (right). In the multiplicity distribution for the VBF case, a large difference arises between generating the third jet through the matrix element or through the parton shower. These deviations make the high-multiplicity values for the VBF process unreliable. However, these uncertainties appear at rather small values for the additional jet multiplicities and, hence, are largely irrelevant phenomenologically.

4 Conclusions

We have analysed several observables that distinguish the two main mechanisms for Higgs production, gluon fusion and vector-boson fusion, by looking at the jet activity and at the final-state
event topology in Higgs + 2-jet events. In particular, we have considered the azimuthal correlation between the two tagging jets and a veto on the jet activity in the rapidity interval between the tagging jets. Our work builds upon previous parton-level work, by adding the parton-shower contribution. We have used ALPGEN to generate the appropriate matrix elements for the primary scattering, and have supplemented it with the parton showers generated by HERWIG.

In the case of the azimuthal correlation, we find that the dip at $\Delta \phi_{jj} = \pi/2$, characteristic of a CP-even Higgs boson produced via gluon fusion, is only slightly filled by the parton shower. It is not permissible, however, to generate also the leading jets by the parton shower. Such a procedure eliminates the characteristic azimuthal angle correlations of gluon fusion events since the shower is generated flat in azimuth. As regards the jet veto, to quantify the jet activity in the rapidity interval between the tagging jets, we have computed the multiplicity distribution of jets within the rapidity interval and normalised it to the total cross section for Higgs + 2-jet production, after jet reconstruction.

It should be cautioned that leading-order calculations for multijet rates, as employed in this contribution, may lead to an unphysical dependence of the jet cross sections on parton-level generation cuts. For example, reducing the minimum $p_T$ for the 3rd parton below the jet identification threshold, e.g. in the Higgs + 3-parton VBF channel, will still allow the generation of events with 3 jets after shower evolution, where the third jet arises from radiation off the initial state or the two final-state hard partons. In absence of the appropriate virtual corrections or Sudakov suppression, the input partonic cross section diverges when the $p_T$ cut is sent to zero, and the rate for these events can become unphysically large.

Acknowledgments

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References

PROSPECTS FOR THE SEARCH FOR HIGGS BOSONS WITH
VECTOR BOSON FUSION PROCESSES AT THE LHC

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The search for the Higgs boson is one of the main physics goals of the Large Hadron Collider (LHC) and its two multi-purpose experiments, ATLAS and CMS. Vector boson fusion is the second largest production process for a standard model Higgs boson at the LHC and offers excellent means for background suppression. This paper gives an overview of the prospects of Higgs boson searches using vector boson fusion at the LHC. For a standard model Higgs boson, the decay channels $H \rightarrow \tau\tau$, $H \rightarrow WW$ and $H \rightarrow \gamma\gamma$ are discussed. The discovery potential in the framework of the MSSM is summarized.

1 The Large Hadron Collider

The Large Hadron Collider (LHC) is a proton-proton collider, which is currently being built at CERN, Geneva. The first physics run at the design center-of-mass energy of 14 TeV is expected for 2008. At first, the LHC will operate at low luminosity ($2 \times 10^{33}$ cm$^{-2}$s$^{-1}$) and will later increase its luminosity to the design value of $10^{34}$ cm$^{-2}$s$^{-1}$. ATLAS$^{1}$ and CMS$^{2}$, the two multi-purpose experiments at the LHC, are designed to investigate a wide range of physics.

2 Vector boson fusion ($qqH$-production)

In the vector boson fusion process, two weak bosons are radiated off the incoming quarks and merge to give the Higgs boson. Vector boson fusion is the second largest production process for a standard model Higgs boson at the LHC.$^{3}$ It has a clear signature which can be used to efficiently suppress the background.

2.1 Tagging Jets

One of the most characteristic features of vector boson fusion processes are the tagging jets that are produced from the quarks that are scattered off the merging heavy vector bosons. These jets typically have high $p_T$ and lie in different hemispheres in the forward- and backward region of the detector.$^{4}$ Very forward jets may not be fully contained in the calorimeter, but partially lie in the beam pipe. Moreover, underlying event$^{5}$ and pile-up$^{6}$ lead to additional energy depositions in the calorimeters. A reliable performance and good understanding of the forward calorimeters as well as efficient energy clustering and jet finding algorithms are therefore needed [Fig. 1].

$^{a}$Additional interactions from the same proton-proton collision as the hard event.
$^{b}$Events from other proton-proton collisions than the hard collision, which might also arise from a later or earlier bunch crossing. For the low luminosity phase, two to three collisions per bunch crossing are expected.
2.2 Central Jet Veto

The decay products of the Higgs boson typically lie in the central detector region. Since there is no colour flow between the quarks in the vector boson fusion process, jet production in the central region is suppressed. In contrast, central emission is favoured in QCD interactions which constitute important background processes at the LHC. A veto on additional jets in the central region is therefore a powerful discriminant between vector boson fusion and QCD background processes such as $t\bar{t}$-production [Fig. 2].

Underlying event and pile-up may produce jets in the central region that do not originate from the vector boson fusion hard interaction. This may considerably alter the performance of the jet veto. A good understanding of underlying event and pile-up levels is crucial for this reason. Because of the large uncertainties of forward jet tagging and the central jet veto in the presence of pile-up, all analyses presented here assume the low luminosity case.

3 Search in the $H \rightarrow \tau\tau$ decay mode

The $H \rightarrow \tau\tau$ channel allows access to a Higgs boson-fermion coupling in the decay. At the LHC, this is in the standard model otherwise only possible in the $H \rightarrow bb$ mode, which is experimentally challenging. For masses of the Higgs boson below $m_H = 140$ GeV, its branching ratio to a $\tau$-lepton pair is around 4-8% but drops at higher masses when the decay to heavy vector bosons opens up. For triggering, at least one high-$p_T$ lepton in the final state is needed. For this

![Figure 3: Reconstructed Higgs boson mass for $H \rightarrow \tau\tau \rightarrow 1\ell$ after analysis cuts at CMS (30 fb$^{-1}$).](image)

![Figure 4: Transverse mass distribution for $H \rightarrow WW \rightarrow e\mu$ after analysis cuts at ATLAS.](image)
reason, the decay chains $H \rightarrow \tau \tau \rightarrow l\nu + 3\nu$ and $H \rightarrow \tau \tau \rightarrow l\nu + 4\nu$ are considered. The main background process is irreducible $Z \rightarrow \tau \tau$ -production from strong and electroweak processes. Also $tt$-production and WW-production in association with jets contribute.

In spite of the neutrinos in the final state, the reconstruction of the Higgs boson mass is possible by use of the collinear approximation. Since the $\tau$-leptons obtain a large Lorentz boost due to the large mass and $p_T$ of the Higgs boson, their decay products are emitted roughly in the direction of the original $\tau$. Exploiting momentum conservation in the transverse plane yields the 4-momentum vectors of the two $\tau$-leptons and thus the Higgs boson invariant mass [Fig. 3]. This channel is promising for a 5$\sigma$-discovery in a mass range of about 120 to 135 GeV with 30fb$^{-1}$.

The Higgs boson will not be observable in this decay mode in inclusive searches.

4 Search in the $H \rightarrow WW$ decay mode

The $H \rightarrow WW$ channel gives clean access to the Higgs boson coupling to W-bosons both in production and in decay. The branching ratio rapidly increases with the Higgs boson mass and reaches values over 96% at masses around $m_H \approx 170$ GeV. It drops to about 70% when the decay to the heavier Z-boson opens up.

The semileptonic $H \rightarrow WW \rightarrow l\nu j j$ and the purely leptonic $H \rightarrow WW \rightarrow l\nu l\nu$ decay modes provide the required lepton for triggering. Especially for the semileptonic mode, large background contributions have to be suppressed. These include $tt$, W, Z, WW and ZZ-production in association with jets as well as QCD multijet production.

Since Higgs boson masses near the production threshold of the W-boson pair are considered, the decay products of the Higgs boson have relatively low $p_T$. The collinear approximation is therefore not usable. Instead, the transverse mass $M_T = \sqrt{2 p_T, l \cdot p_T, \text{miss}(1 - \cos \Delta \phi)}$ is calculated [Fig. 4].

In case of the $H \rightarrow WW \rightarrow l\nu j j$ -decay, the Higgs boson mass can also be reconstructed by using the known W-boson mass as an input.

With a cut analysis method, this channel is promising for a 5$\sigma$-discovery in a mass range of 125 to 190 GeV with 30fb$^{-1}$. Only 5fb$^{-1}$ are needed in the mass region of 150 to 190 GeV.

5 Search in the $H \rightarrow \gamma\gamma$ decay mode

The branching ratio of the standard model Higgs boson to two photons is at most 0.22% at $m_H \approx 125$ GeV. However, this decay channel has an excellent Higgs boson mass resolution. The main background is photon production in association with jets. Processes with final state electrons and jets contribute if these particles are misidentified as isolated photons.

To achieve a final mass resolution of $\sigma_{\gamma\gamma} = 0.7\%$, it is necessary to reconstruct the Higgs boson vertex to reduce the photon direction uncertainty. An efficient method for vertex reconstruction is to determine the longitudinal position of the highest $p_T$ track in the event.

Using a cut analysis method, a significance of 2.5$\sigma$ can be achieved for a Higgs boson mass of 120 GeV with 30fb$^{-1}$. As a discovery channel, vector boson fusion $H \rightarrow \gamma\gamma$ is not competitive with the inclusive mode which reaches significances above 5$\sigma$.

6 Higgs bosons of the Minimal Supersymmetric Standard Model (MSSM)

The MSSM features two SU(2) Higgs doublets with a total of five observable states, two neutral CP-even, one neutral CP-odd and one charged Higgs boson. The Higgs sector can be described at tree level by two parameters, which are usually chosen to be the mass of the CP-odd Higgs boson, $M_A$, and the ratio of the vacuum expectation values of the two Higgs doublets, $\tan \beta = \frac{\langle v_1 \rangle}{\langle v_2 \rangle}$.

Only the CP-even Higgs bosons can be produced in vector boson fusion. At 30fb$^{-1}$, at least one of these Higgs bosons will be observable over the entire $M_A$-$\tan \beta$ plane in the $m_H^{\text{max}}$, no-mixing,
small $\alpha_{	ext{QCD}}$ and the gluophobic benchmark scenario\cite{13,14}. Inclusive searches need higher integrated luminosities for complete coverage of the $M_A$-$\tan \beta$ plane\cite{13}.

7 Conclusions

Vector boson fusion offers an outstanding discovery potential for the standard model Higgs boson, especially in the low mass region between the LEP limit of 114.4 GeV\cite{15} and 190 GeV [Fig. 5, 6]. In the supersymmetric framework of the MSSM, vector boson fusion processes have an excellent discovery potential over the whole parameter plane.

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References

INCLUSIVE HIGGS BOSON SEARCHES IN FOUR-LEPTON FINAL STATES
AT THE LHC

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The inclusive search for the Standard Model Higgs boson in four-lepton final states with the
ATLAS and CMS detectors at the LHC pp collider is presented. The discussion focuses
on the $H \rightarrow ZZ^{(*)} \rightarrow 4l + X$ decay mode for a Higgs boson in the mass range $120 \leq M_H \leq 600$ GeV/$c^2$. A prospective analysis is presented for the discovery potential based on a
detailed simulation of the detector response in the experimental conditions of the first years of
LHC running at low luminosity. An overview of the expected sensitivity in the measurement
of the Higgs boson properties is also given.

1 Introduction

The Standard Model(SM) of electroweak interactions contains one Higgs boson whose mass,
$M_H$, is a free parameter of the model. The inclusive single production reaction $p + p \rightarrow H + X$
followed by the decay $H \rightarrow ZZ^{(*)} \rightarrow l^+l^-l'^+l'^-$ (in short $H \rightarrow 4l$) is the clearest("golden") decay
mode for the discovery of the SM Higgs boson at the LHC and can provide a sensitivity over a
wide range of masses $M_H$ from 120 to 600 GeV/$c^2$. There are three different final states which
depend on the flavour of the Z-boson decay leptons: $H \rightarrow 4e$, $H \rightarrow 4\mu$ and $H \rightarrow 2e2\mu$. Thanks
to the relatively small background contamination, the $H \rightarrow 4l$ also allows a precise measurement
of the Higgs boson properties (mass, width, spin, couplings, etc...).

The report summarizes the expected potential of ATLAS and CMS in SM Higgs boson
searches in the $H \rightarrow 4l$ channel. For more details on the analyses described here, the reader is
directed to the ATLAS$^1$ and CMS$^2$ Physics Technical Design Reports.

*aNow at the Universiteit Antwerpen.*
2 Higgs boson signal and backgrounds

At LHC energies, there are two dominant SM Higgs boson production processes: gluon-gluon, \( gg \to H \), and weak boson fusion \( qq \to q\bar{q}H \). In the mass range \( M_H \lesssim 135 \text{ GeV}/c^2 \), the SM Higgs boson decays mainly into \( bb \) (BR \( \approx 85\% \)) and \( \tau^+\tau^- \) (BR \( \approx 8\% \)) pairs but the search in the H \( \to \gamma\gamma \) decay mode is privileged despite it small branching ratio (BR \( \approx 0.2\% \)) because of its clean experimental signature. For \( M_H \gtrsim 135 \text{ GeV}/c^2 \), the decay into H \( \to W^+W^- \) is dominant. The branching ratio for the H \( \to ZZ^{(*)} \) decay is sizable for \( M_H \geq 115 \text{ GeV}/c^2 \).

The H \( \to ZZ^{(*)} \to 4l \) signal event topology is characterized by two pairs of oppositely charged and isolated same-flavour leptons coming from the same vertex with a di-lepton invariant mass compatible with the Z-boson mass. The Higgs boson signal manifests itself as a narrow mass peak in the reconstructed four-lepton invariant mass spectrum. There are three main background sources to the H \( \to 4l \) signal. The reducible Zbb and \( tt \) background processes differ from the signal by the presence of two non-isolated leptons inside b-jets with displaced vertices. The irreducible ZZ\((*)\) background has kinemactical properties which are very similar to that of the signal except for the four-lepton invariant mass which shows a broad spectrum.

3 Event selection

All selections are optimised to have the highest significance for discovery with emphasis on a realistic strategy for the control of experimental errors and background systematics.

The on-line preselection consists of a logical OR of basic single and double electron or muon triggers. The off-line preselection starts with the search for events with at least four lepton candidates within the fiducial volume. The aim is to reduce as much as possible the contamination of background sources involving "fake" leptons from QCD jets while preserving as much as possible the signal detection efficiency. The Zbb and \( tt \) backgrounds have at least one non-isolated lepton-pair with often detectable displaced vertices in contrast to the signal and ZZ\((*)\) background. Therefore, the most discriminating preselection variables against these backgrounds come from vertex constraints and isolation criteria relying on the measurement of primary tracks in the tracker and/or the energy flow in the calorimeters.

The kinematical selection consists of cuts on the lepton transverse momenta and the reconstructed di-lepton invariant mass spectra. The first cut exploits the fact that b-decay leptons from the Zbb and \( tt \) backgrounds have on average a softer \( p_T \) spectrum than leptons from the Higgs boson signal or ZZ\((*)\) background. The second requirement is powerful against all backgrounds.

The number of signal and background events is determined by a simple window sliding in the reconstructed four-lepton invariant mass spectrum. After the full selection, the reducible backgrounds are suppressed well below the level of the ZZ\((*)\) contamination which remains the dominant and sole remaining background. The typical rejection factors vary from \( 2 \times 10^3 \) to \( 10^4 \) for \( tt \), from \( 500 \) to \( 10^5 \) for Zbb and from \( 20 \) to \( 4 \) for ZZ\((*)\), depending on the M_H-hypothesis, for a signal selection efficiency of 25-55 \%.

4 Systematics

The systematics on the signal significance are related to the knowledge of the ZZ\((*)\) background rate in the signal region and the uncertainty on this knowledge. Two approaches have been followed to estimate the background directly from the data: a normalisation to single Z, Z \( \to 2l \), data and a normalisation to sidebands. Both approaches lead to a reduced sensitivity to theoretical and experimental uncertainties as well as a full cancellation of the luminosity uncertainty. The theoretical uncertainty is of the order of 2 to 8\% for the normalisation to Z \( \to 2l \) and 0.5
Figure 1: The expected statistical significance for the Standard Model Higgs boson signal as function of its mass for an integrated luminosity of 30 fb$^{-1}$ for the ATLAS (Left) and CMS (Right) experiments.

to 4% for the normalisation to sidebands. The low statistics of ZZ(*) events could be a limiting factor for the sidebands method.

The overall strategy for controlling the detector systematics is to estimate the efficiency and the precision of the energy and momentum measurements from experimental data. Single Z and single W processes have huge cross sections at the LHC, and are expected to lead to a significant reduction of the reconstruction uncertainties already after few fb$^{-1}$.

5 Discovery reach

Figure 1 shows the expected statistical significance for the SM Higgs boson signal as function of its mass for an integrated luminosity of 30 fb$^{-1}$. A 5$\sigma$-discovery is possible over a wide range of masses in the $H \to 4l$ channel: $130 < M_H < 160$ GeV/c$^2$ and $2m_Z < M_H < 550$ GeV/c$^2$. The drop in sensitivity around $M_H \approx 180$ GeV/c$^2$ will be filled by the complementary channel $H \to WW^{(*)} \to 2\ell 2\nu$ where less than 1 fb$^{-1}$ is needed for 5$\sigma$-discovery. For $M_H < 130$ GeV/c$^2$, the highest discovery potential is obtained in the $H \to \gamma \gamma$ decay mode.

6 Measurement of the Higgs boson properties

The Higgs boson mass and width are obtained from a fit to the four-lepton invariant mass spectrum. For an integrated luminosity of 30 fb$^{-1}$, the expected statistical precision of the mass measurement is better than 1% over a wide range of masses (Figure 2 Left). The Higgs boson width measurement is only possible for Higgs boson masses beyond 200 GeV/c$^2$ when the Higgs boson natural width starts to dominate the experimental resolution. The expected precision on the width is smaller than 30%.

The Higgs boson couplings to fermions and gauge bosons can be extracted from rate measurements in the different Higgs boson production and decay channels. Relative precision on the squared Higgs boson couplings, assuming 300 fb$^{-1}$ of data collected by both the ATLAS and the CMS experiments, varies between 10% and 40% depending on the coupling, except for the Yukawa coupling to the b-quark which suffers from large uncertainties related to b-tagging and background normalisation.
The $H \rightarrow 4l$ channel is particularly suitable to measure the Higgs boson spin and CP state because its small background contamination and the fact that the event kinematics can be completely reconstructed with good precision. The angular correlations between the $Z$-boson decay products are used to extract the spin and CP state of the resonance. A study based on ATLAS fast simulation shows that with 100 fb$^{-1}$ a pseudo-scalar Higgs boson can be ruled out if $M_H > 200$ GeV/c$^2$ and an axial vector and vector Higgs boson can be excluded if $M_H > 230$ GeV/c$^2$. A recent analysis by the CMS experiment considers also CP-violating spin-0 Higgs boson states via the introduction of a CP-mixing parameter and determines the minimal enhancement or suppression in cross section needed in order to exclude the SM pseudo-scalar Higgs boson. It is shown that the distinction between a scalar and a pseudo-scalar Higgs boson is already possible with 60 fb$^{-1}$ integrated luminosity.

Acknowledgments

The author would like to thank M.Aldaya, S.Baffioni, F. Beaudette, C.Charlot, A.Drozdetskiy, D.Fuytan, A.Nikitenko and I.Puljak for many interesting and helpful discussions.

References

6. M.Bluj, A study of angular correlations in $\phi \rightarrow ZZ \rightarrow 2e2\mu$, CMS Note 2006/094.
Evidence for Electroweak Single Top Quark Production at D0

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We present first evidence for electroweak single top quark production from the D0 Collaboration at the Fermilab Tevatron Collider. Using a 0.9 fb⁻¹ data sample, several multivariate techniques are used to isolate the single top signal from background. Combining these three methods (Bayesian Neural Networks, Decision Trees, and Matrix Elements), we obtain a measured cross section of \( \sigma(p\bar{p} \rightarrow t\bar{b} + X, t\bar{q}b + X) = 4.8 \pm 1.3 \text{ pb} \) with a Gaussian significance of 3.5σ. Using this measurement, we set a lower 95% C.L. on the \( V_{tb} \) element of the Cabibbo-Kobayashi-Maskawa (CKM) matrix of \( 0.68 < |V_{tb}| \leq 1 \).

1 The Single Top Analysis

1.1 Theory and Motivation

Examining single top quark production at the Tevatron¹, we find two major production modes for single top and Figure 1 shows the two representative Feynman diagrams. The first one is called s-channel production where two quarks annihilate to make an off-shell \( W \) boson which then decays to a top quark and a bottom quark. This process is also known as \( t\bar{b} \) because it has a top and bottom quark in the final state. The second process, or \( b \)-channel production, is when a light quark emits a \( W \) boson which fuses with a \( b \)-quark coming from a gluon that splits into a \( t\bar{b} \) pair to make a top quark. Because this process has a top, a bottom, and a light quark in the final state, it is also known as \( t\bar{q}b \).

Observation of single top quark production allows one to study the \( V_{tb} \) coupling and directly measure the CKM element \( V_{tb} \) without the standard model (SM) assumption of three families. In addition, the single top cross section is also sensitive to new physics. Since single top is a third process called associated production, but the cross section is too small to observe at the Tevatron, so we will not discuss it further.

¹There is a third process called associated production, but the cross section is too small to observe at the Tevatron, so we will not discuss it further.
Figure 1: Representative Feynman diagrams for single top quark production. The left plot shows the $t\bar{b}$ production mode and the right plot shows the $t\bar{g}b$ production mode.

A background for the Higgs search, this analysis is a “proof of principle” for the advanced techniques that are being used.

1.2 Experimental Signature, Backgrounds, and Data Samples

The experimental signature for single top is a high $p_T$ lepton, large missing transverse energy corresponding to the neutrino, and two to four jets (this is done to maximize acceptance). One can further enhance the signal content of the sample by requiring that either one or two jets are associated with $b$-quarks (also known as $b$-tagging).

There are three main backgrounds to single top. One comes from $\bar{t}t$ pair production which is normalized to the NNLO cross section. Then there is multijet production where a jet fluctuates to mimic a lepton which also produces missing energy. And $W+$ jet production which includes $W\bar{t}\bar{t}, W\bar{c}c,$ and $Wjj$. The multijet and $W+$ jet background yields are normalized to data before $b$-tagging.

We take our full data sample and divide it into electrons and muons. These samples are further separated by the number of jets in the event. Finally, we divide these samples again by determining how many jets can be associated with $b$-quark(s). The event yields and systematic uncertainties are shown in Table 1. The acceptance for the single top signal is $(3.2 \pm 0.4)\%$ for $tb$ and $(2.1 \pm 0.3)\%$ for $tqb$. In addition, we have checked the agreement between data and background model for over 100 individual distributions and find good agreement.

<table>
<thead>
<tr>
<th>Source</th>
<th>2 jets</th>
<th>3 jets</th>
<th>4 jets</th>
</tr>
</thead>
<tbody>
<tr>
<td>$tb$</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>$\bar{t}t \rightarrow \ell\ell$</td>
<td>39±9</td>
<td>32±7</td>
<td>11±3</td>
</tr>
<tr>
<td>$\bar{t}t \rightarrow \ell+jets$</td>
<td>20±5</td>
<td>103±25</td>
<td>143±33</td>
</tr>
<tr>
<td>$W\bar{b}\bar{b}$</td>
<td>261±55</td>
<td>120±24</td>
<td>35±7</td>
</tr>
<tr>
<td>$W\bar{c}c$</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 bkgd</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>

Table 1: The left table show the numbers of expected and observed events in 0.9 $fb^{-1}$ for $e$ and $\mu$, 1 $b$ tag and 2 $b$ tag channels combined. The right table shows the sizes of the systematic uncertainties for this analysis.

*An important point to note is that the single top signal is smaller than the background uncertainty.
Figure 2: The top left plots shows the full discriminant for the Matrix Element Analysis. The top right plots shows an enlarged region of the discriminant near the value one. The bottom left plot shows an enlarged region of the Decision Tree output for high values of the discriminant. The bottom right plot shows the posterior distributions for the expected and measured cross sections for the Decision Tree analysis.

<table>
<thead>
<tr>
<th></th>
<th>Bayesian NN</th>
<th>Matrix Element</th>
<th>Decision Trees</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sigma(p^+ p^- \rightarrow t\bar{b} + X, t\bar{q}b + X)$ [pb]</td>
<td>$3.2^{+1.3}_{-1.8}$</td>
<td>$5.0 \pm 1.9$</td>
<td>$3.0^{+1.5}_{-1.3}$</td>
</tr>
<tr>
<td>Significance</td>
<td>$1.3\sigma$</td>
<td>$2.4\sigma$</td>
<td>$1.8\sigma$</td>
</tr>
</tbody>
</table>

Table 2: Table of expected and observed cross sections along with the expected sensitivities and observed significances.

1.3 Multivariate Methods and Results

We have used three methods to separate the single top signal from the background. The first is called Bayesian Neural Networks, the second is called Decision Trees, and the third is called Matrix Elements. Then using the outputs from each method we construct a binned likelihood and extract a cross section.

The first method uses Bayesian Neural Networks. It uses 24 input variables for training the networks (where the signal tends towards one and the background towards zero). A simple description of BNNs is that of averaging many individual NNs. The second method uses Decision Trees. This analysis uses 49 input variables, and the idea is to recover events that are rejected by a simple cut-based analysis. In addition, for events that get misclassified, we use boosting which effectively averages over many trees to improve signal and background separation. The third method used is called the Matrix Element Method. This is a different idea based on using the full kinematic information from the reconstructed objects in the event to form a discriminant to separate signal from background.

In Figure 2, we show the discriminant output for the Matrix Element and Decision Tree analyses along with the expected and measured cross sections from the Decision Tree analysis. In Table 1.3, we show the expected sensitivities and the measured values of the cross sections along with the significances. Decision trees have a 3.4 excess which establishes the first evidence for single top quark production.
The most general $tbW$ vertex is given by\(^5\):

\[
\Gamma^W_{tbW} = -\frac{g}{\sqrt{2}} V_{tb} \bar{u}(p_b) \left[ g^\mu (f_1^L P_L + f_1^R P_R) - \frac{i g^\mu W}{M_W} (f_2^L P_L + f_2^R P_R) \right] u(p_t),
\]

where $k$ is the $W$ four-momentum and the $f_1$ and $f_2$ couplings can a-priori be CP-violating. In the case of the SM, CP is conserved in the $tbW$ vertex and $f_1^L = 1$ and $f_1^R = f_2^L = f_2^R = 0$. Since Decision Trees are the most sensitive result, with a measured cross section of $\sigma(pp \to tb + X, tgb + X) = 4.9 \pm 1.4 \text{ pb}$, we use this value to extract a value of $|V_{tb} f_1^L| = 1.3 \pm 0.2$ and assuming full SM values (i.e. $f_1^L = 1$) we set a lower limit on $|V_{tb}|$ at 95% C.L. of $0.68 < |V_{tb}| \leq 1$. We have combined the outputs of these three measurements using a simple linear method and find a combined cross section of $4.8 \pm 1.3 \text{ pb}$ with a significance of $3.5 \sigma$. We show these results in Figure 3.

In summary, we have found first evidence for single top quark production with a measured cross section $\sigma(pp \to tb + X, tgb + X) = 4.8 \pm 1.3 \text{ pb}$. Using this measurement we have set lower limits on $|V_{tb}|$ of $0.68 < |V_{tb}| \leq 1$.

Acknowledgments

We thank the staffs at Fermilab and collaborating institutions, and acknowledge support from the DOE and NSF (USA); CEA and CNRS/IN2P3 (France); FASI, Rosatom and RFBR (Russia); CAPES, CNPq, FAPERJ, FAPESP and FUNDUNESP (Brazil); DAE and DST (India); Colciencias (Colombia); CONACyT (Mexico); KRF and KOSEF (Korea); CONICET and UBACyT (Argentina); FOM (The Netherlands); PPARC (United Kingdom); MSMT (Czech Republic); CRC Program, CPT, NSERC and WestGrid Project (Canada); BMBF and DFG (Germany); SFI (Ireland); The Swedish Research Council (Sweden); Research Corporation; Alexander von Humboldt Foundation; and the Marie Curie Program.

References

SEARCH FOR SINGLE-TOPO PRODUCTION AT CDF

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This article reports on recent searches for single-top-quark production by the CDF collaboration at the Tevatron using a data set that corresponds to an integrated luminosity of 955 pb$^{-1}$. Three different analyses techniques are employed, one using likelihood discriminants, one neural networks and one matrix elements. The sensitivity to single-top production at the rate predicted by the standard model ranges from 2.1 to 2.6σ. While the first two analyses observe a deficit of single-top like events compared to the expectation, the matrix element method observes an excess corresponding to a background fluctuation of 2.3σ. The null results of the likelihood and neural network analyses translate in upper limits on the cross section of 2.8 pb for the t-channel production mode and 3.7 pb for the s-channel mode at the 95% C.L. The matrix element result corresponds to a measurement of 2.7$^{+1.5}_{-1.3}$ pb for the combined t- and s-channel single-top cross section.

In addition, CDF has searched for non-standard model production of single-top-quarks via the s-channel exchange of a heavy W' boson. No signal of this process is found resulting in lower mass limits of 780 GeV/c$^2$ in case the mass of the right-handed neutrino is smaller than the mass of the right-handed W' or 790 GeV/c$^2$ in the opposite case.

Keywords: single-top; electroweak top quark production; W' boson.

1 Introduction

In pp collisions at the Tevatron top quarks are mainly produced in pairs via the strong force. However, the standard model also predicts the production of single top-quarks by the weak interaction via the s- or t-channel exchange of an off-shell W boson. While early Run II searches by the CDF and DØ collaborations, based on data sets corresponding to 162 pb$^{-1}$, 230 pb$^{-1}$ or 695 pb$^{-1}$ of integrated luminosity, did not find evidence for single-top production$^{1,2,3}$, one of the latest DØ analyses$^4$ using data with $L_{int} = 1$ fb$^{-1}$ observes an excess of single-top-like events of 3.4σ.

The single-top production cross section is predicted to $\sigma_{t+}$ = 2.9 $\pm$ 0.4 pb for a top mass of 175 GeV/c$^2$ which is about 40% of the top-antitop pair production cross section. The main obstacle in finding single-top is however not the production rate of the signal but the large background rate. After all selection requirements are imposed, the signal to background ratio is approximately 1/16. This challenging, background-dominated dataset is the main motivation for using multivariate techniques.

2 Standard Model Searches

In this article we present three new CDF searches for standard model single-top production and one new search for a W' boson decaying into tb. All four analyses are based on the same event selection and use the same Run II data set corresponding to an integrated luminosity of 953 pb$^{-1}$. The event selection exploits the kinematic features of the signal final state, which contains a top quark, a bottom quark,
Table 1: Expected number of signal and background events and total number of events observed in 955 pb\(^{-1}\) in the CDF single-top dataset.

<table>
<thead>
<tr>
<th>Process</th>
<th>(N_{\text{events}})</th>
<th>Process</th>
<th>(N_{\text{events}})</th>
</tr>
</thead>
<tbody>
<tr>
<td>(W + bb)</td>
<td>170.9 ± 50.7</td>
<td>non-(W)</td>
<td>26.2 ± 15.9</td>
</tr>
<tr>
<td>(W + c\bar{c})</td>
<td>63.5 ± 19.9</td>
<td>(t\bar{t})</td>
<td>58.4 ± 13.5</td>
</tr>
<tr>
<td>(Wc)</td>
<td>68.6 ± 19.8</td>
<td>Diboson</td>
<td>13.7 ± 1.9</td>
</tr>
<tr>
<td>Mistags</td>
<td>136.1 ± 19.7</td>
<td>(Z + \text{jets})</td>
<td>11.9 ± 4.4</td>
</tr>
<tr>
<td>Total background</td>
<td>549.3 ± 95.2</td>
<td></td>
<td></td>
</tr>
<tr>
<td>(t)-channel</td>
<td>22.4 ± 3.6</td>
<td>(s)-channel</td>
<td>15.4 ± 2.2</td>
</tr>
<tr>
<td>Total prediction</td>
<td>587.1 ± 96.6</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Observed</td>
<td>644</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

and possibly additional light quark jets. To reduce multi-jet backgrounds, the \(W\) originating from the top quark is required to have decayed leptonically. One therefore demands a single, isolated high-energy electron or muon \((E_T(e) > 20 \text{ GeV}, \text{ or } E_T(\mu) > 20 \text{ GeV}/c)\) and large missing transverse energy from the undetected neutrino \(E_T > 25 \text{ GeV}\). Electrons are measured in the central and in the forward calorimeter, \(|\eta| < 2.0\). To further suppress events in which no real \(W\) boson is produced, called non-\(W\) background, additional cuts are applied. The cuts are based on the assumption that these events do not produce \(E_T\) by nature but due to lost or mismeasured jets. Therefore, one would expect small \(E_T\) and small values of the angle \(\Delta \phi\) between \(E_T\) and a jet. We further reject dilepton events from \(Z\) decays by requiring the dilepton mass to be outside the range: \(76 \text{ GeV}/c^2 < M_{\mu\mu} < 106 \text{ GeV}/c^2\).

The remaining backgrounds belong to the following categories: \(Wbb\), \(Wc\bar{c}\), \(Wc\), mistags (light quarks misidentified as heavy flavor jets), non-\(W\) and diboson \(WW\), \(WZ\), and \(ZZ\). We remove a large fraction of the backgrounds by demanding exactly two jets with \(E_T > 15 \text{ GeV}\) and \(|\eta| < 2.8\) be present in the event. At least one of these two jets should be tagged as a \(b\) quark jet by using displaced vertex information from the silicon vertex detector (SVX). The numbers of expected and observed events are listed in table 1.

2.1 Likelihood Discriminant Analysis

One multi-variate analysis uses likelihood discriminants to combine several variables to a discriminant to separate single-top events from background events. One likelihood discriminant is defined for the \(t\)-channel, one for the \(s\)-channel search. Seven or six variables are used, respectively. The likelihood functions are constructed by first forming histograms of each variable. The histograms are produced separately for signal and several background processes. The histograms are normalized such that the sum of their bin contents equals 1. For one variable the different processes are combined by computing the ratio of signal and the sum of the background histograms. These ratios are multiplicatively combined to form the likelihood discriminant.

One of the variables used in the analysis is the output of a neural net \(b\) tagger. In figure 1a the distribution of this \(b\) tag variable is shown for the 644 data events. In case of double-tagged events the leading \(b\) jet (highest in \(E_T\)) is included in this distribution. The neural net \(b\) tagger gives an additional handle to reduce the large background components where no real \(b\) quarks are contained, mistags and charm-backgrounds. Both of them amount to about 30% in the \(W + 2\) jets data sample even after imposing the requirement that one jet is identified by the secondary vertex tagger of CDF\(^8\).

The \(t\)-channel likelihood function is shown in figure 1b. The best sensitivity (expected p-value 2.5%) is reached by combining the two likelihood discriminants in a two-dimensional fit where \(t\)-channel and \(s\)-channel are considered as one single-top signal (combined search). The observed data show no indication of a single-top signal and are compatible with a background-only hypothesis (p-value 38.5%). The upper limit on the combined single-top cross section is found to be 2.7 pb at the 95% C.L., while the expected limit is 2.9 pb. The best fit for the cross sections yields \(\sigma_t = 0.2^{+0.9}_{-0.2}\) pb and \(\sigma_s = 0.1^{+0.7}_{-0.1}\) pb.

2.2 Neural Network Search

In the second analysis a neural network is used to combine 23 kinematic or event shape variables are combined to a powerful discriminant. Figure 2a shows the observed data for the combined search compared
to the expectation in the signal region defined by a neural network output between 0.3 and 1. In the combined search where the ratio of $t$-channel and $s$-channel cross sections is fixed to the standard model value a p-value of 54.6% is observed, providing no evidence for single-top production. The corresponding upper limit on the cross section is 2.6 pb at the 95% C.L. To separate $t$- and $s$-channel production two additional networks are trained and a simultaneous fit to both discriminants is performed. The best fit values are $\sigma_t = 0.5^{+1.1}_{-0.2}$ pb for the $t$-channel and $\sigma_s = 0.7^{+1.6}_{-0.7}$ pb. The corresponding upper limits are 2.6 pb and 3.7 pb, respectively. The observed p-value is 21.9%.

### 2.3 Matrix Element Analysis

Another way to discriminate signal from background is to compute leading order matrix elements for signal and background processes. The measured four-vectors of the jets and the charged lepton are used as experimental input. Constraints on energy and momentum conservation are applied. The jet energy measurements are corrected to parton level energies using transfer functions. To obtain a relative weight for a certain hypothesis one integrates over jet energies and the momenta of the incoming partons using the parton density functions as integration kernel. The weights for the individual hypotheses are combined event-by-event by forming a ratio signal and signal-plus-background weights. The resulting discriminant is named event probability density (EPD). The measured EPD distribution is shown in figure 3, compared to the fitted event rates for the different processes. A single-top signal corresponding to a 2.3 $\sigma$ excess is observed. The associated single-top cross section is $2.7^{+1.6}_{-1.3}$ pb.
3 Search for a $W'$ boson

Based on a the same event selection as the standard model searches CDF has also searched for a $W'$ boson in the decay channel, $W' \rightarrow tb$. The signal is modeled using the event generator PYTHIA \textsuperscript{7}. The invariant mass of the charged lepton, the reconstructed neutrino and the two leading jets is used as a discriminant. The measured data are shown in comparison to the standard model prediction in figure 3b. No evidence for a resonant $W'$ boson production is found, yielding limits on the production cross section ranging from $2.3 \text{ pb}$ at $M(W') = 300 \text{ GeV}$ to $0.4 \text{ pb}$ at $M(W') = 950 \text{ GeV}$. Utilizing theoretical cross section calculations\textsuperscript{8}, lower limits on the $W'$ mass are set: $M(W') > 760 \text{ GeV}$ if the mass of potential right-handed neutrinos is below $M(W')$ and $M(W') > 790 \text{ GeV}$ otherwise.

4 Conclusions

The collaboration has performed searches for standard and non-standard model single-top production using a data set corresponding to an integrated luminosity of $955 \text{ pb}^{-1}$. No evidence for these processes could be established. Two standard model searches based on likelihood discriminants or neural networks find no excess which can be attributed to single-top production, while the matrix element analysis finds an excess of $2.3\sigma$ compatible with single-top production. The overall consistency of all these analyses is only 1%, since the analyses feature a correlation between 60% and 70%. At present, there are no hints to other causes than statistical fluctuations. Even larger data samples will be available in the near future and will help clarify the situation. The $W'$ search in the single-top channel establishes new upper limits of $M(W') > 760 \text{ GeV}$ or $M(W') > 790 \text{ GeV}$ depending on the mass of the right-handed neutrino.

References

4. V. M. Abazov et al. (DO Collaboration), hep-ex/0612052.
QCD and High Energy Hadronic Interactions

VIII - Top and W Session

Chairperson: B. Bloch-Devaux
W mass and width measurements at the Tevatron

Emily Nurse (for the CDF collaboration)

University College London, Gower Street, London, WC1E 6BT, United Kingdom.

I present a measurement of the W boson mass ($M_W$) and width ($\Gamma_W$) using 290 and 350 pb$^{-1}$ of CDF Run II data respectively. The measurements, performed in both the electron and muon decay channels, rely on a fit to the W transverse mass distribution. We measure $M_W = 80413 \pm 48$ MeV and $\Gamma_W = 2032 \pm 71$ MeV which represent the world's single most precise measurements to date.

1 Introduction

The mass ($M_W$) and width ($\Gamma_W$) of the W boson are important parameters of the Standard Model (SM). Radiative corrections to the W propagator are dominated by Higgs and top-bottom loops, thus a precise measurement of $M_W$ together with $M_t$, the mass of the top quark, place an indirect constraint on the mass of the as yet un-discovered Higgs boson, $M_H$. A precise measurement of $\Gamma_W$ provides a stringent test of the SM prediction which is accurate to 2 MeV$^1$.

At the Tevatron W bosons are predominantly produced via quark anti-quark annihilation. The measurements are performed in the $e\nu$ and $\mu\nu$ decay channels which provide clean experimental signatures. The $M_W$ and $\Gamma_W$ analyses utilise 200 pb$^{-1}$ and 350 pb$^{-1}$ of CDF data from Run II at the Tevatron respectively.

Since neutrinos are not detected in CDF the W invariant mass cannot be reconstructed. Instead we reconstruct the transverse mass, $M_T$, which is defined as:

$$M_T = \sqrt{2p_T^e p_T^\nu (1 - \cos \phi_{e\nu})}$$

(1)

where $p_T^e$ is the transverse momentum ($p_T$) of the charged lepton, $p_T^\nu$ is the $p_T$ of the neutrino and $\phi_{e\nu}$ is the azimuthal angle between the charged lepton and the neutrino. $p_T^e$ is inferred from the transverse momentum imbalance in the event.
A Monte Carlo simulation is used to predict the $M_T$ distribution as a function of $M_W$ and $\Gamma_W$. These predictions are fitted to the data with a binned maximum-likelihood fit in order to extract $M_W$ and $\Gamma_W$. The fit for $M_W$ is performed in the region around the peak of the distribution: 65–90 GeV. The fit for $\Gamma_W$ is performed in the high $M_T$ tail region: 90–200 GeV, which is still sensitive to the Breit-Wigner line-shape but less sensitive to the Gaussian detector resolutions. These line-shape predictions depend on a number of production and detector effects. The most important effects are described in this document and all the systematic uncertainties are summarised at the end.

2 Monte Carlo Simulation

A dedicated parameterised Monte Carlo simulation is used to generate the $M_T$ templates used in the fits. The W $p_T$ spectrum is modelled with RESBOS$^2$ and QED corrections for one photon emission are simulated with Berends and Kleiss$^3$ and WGRAD$^4$. Systematic uncertainties arise from non-perturbative QCD parameters affecting the W $p_T$ spectrum and considerations of the emission of a second photon from the final state charged lepton. Parton distribution functions (PDFs) affect the acceptance and kinematics of decay products. The templates are generated with the CTEQ6M$^5$ PDFs and their error sets are used to estimate the PDF uncertainty.

The detector response model is tuned to $Z \rightarrow \ell\ell$ and $W \rightarrow \ell\nu$ data as well as a full GEANT Monte Carlo simulation of the CDF detector.

3 Lepton Calibration: Scales and Resolutions

The muon momentum is measured in a cylindrical drift chamber. The scale and resolution of the momentum are calibrated using the resonance peaks in $J/\Psi \rightarrow \mu\mu$, $\Upsilon(1S) \rightarrow \mu\mu$ and $Z \rightarrow \mu\mu$ events utilising the precisely measured world average masses of these particles.$^1$ The $J/\Psi$ sample has sufficient statistics to verify the linearity of the momentum scale by studying its variance as a function of muon $p_T$. Combining all three measurements enables an accuracy of 0.021% on the momentum scale.

The electron energy is measured in the calorimeter. The electron momentum is also measured in the drift chamber$^6$, thus the well calibrated momentum measurement is used to calibrate the calorimeter scale (response) and resolution using the ratio between the electron energy measured in the calorimeter and the track momentum ($E/p$) in $W \rightarrow e\nu$ events. The scale and resolution can also be obtained independently from the mass peak in $Z \rightarrow ee$ events. The two measurements are combined to give a calorimeter scale measurement accurate to 0.034%.

4 Hadronic Recoil Calibration

The neutrino $p_T$ is determined from the missing transverse energy, $E_T$, in the detector. A recoil vector, $\vec{U}$, is defined as the vector sum of transverse energy over all calorimeter towers, excluding those surrounding the lepton. The $E_T$ is then defined as $-(\vec{U} + \vec{p_T})$. The recoil has contributions from initial state gluon radiation from the incoming quarks, underlying event energy and final state photon radiation from the charged lepton. The recoil is represented by a parameterised model, which is tuned in $Z \rightarrow \ell\ell$ events. The model parameters are found from the Z data and applied to the W data. The systematic uncertainties on $M_W$ and $\Gamma_W$ come from the uncertainties on the model parameters due to the limited statistics in the Z data.

$^6$Since the mass of the electron is negligible the true momentum and energy values are the same. However collinear photon radiation from the electron which is clustered back into the energy measurement can decrease the track momentum measurement. These effects are well modelled in the simulation.
Table 1: Uncertainties for the $W$ mass (left) and $W$ width (right). The third column lists the uncertainties that are common between the electron and muon channels.

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5 Backgrounds

Backgrounds have different $M_T$ distributions to $W \to \ell\nu$ events, therefore the $M_T$ shape must be added to the Monte Carlo templates when fitting to the data. Electroweak backgrounds consist of $Z \to \ell\ell$ events where one of the leptons goes undetected and $W \to \tau\nu$ and $Z \to \tau\tau$ events where the $\tau$ decays to an electron or muon. These backgrounds are found using Pythia 6 Monte Carlo samples of $W$ and $Z$ events, passed through a full GEANT simulation of the CDF detector. Non-electroweak backgrounds consist of multi-jet events, where one jet fakes or contains a lepton and the other is sufficiently mis-measured to produce $E_T$, and (in the muon channel only) kaons that decay to muons within the volume of the drift chamber. In the latter case the resulting reconstructed track contains a kink that can produce a fake high measured $p_T$ and $E_T$. The multi-jet background normalisation is found by fitting the low $E_T$ distribution where this background dominates. The $M_T$ distributions are found by reversing certain lepton identification cuts. The kaon background is found by fitting the high tail of the track fit $\chi^2$ distribution where this background is large. The $M_T$ shape is found by reversing an impact parameter cut.

6 Results

The systematic and statistical uncertainties for $M_W$ and $\Gamma_W$ are summarised in Table 1. Figure 1 shows the $M_T$ fits for $M_W$ in the muon and electron decay channels. The fitted $M_W$ values are combined together with fits to the charged lepton $p_T$ and $E_T$ distributions to give $M_W = 80413 \pm 48$ MeV, the world's most precise single measurement. This result increases the world average central value by 6 MeV and reduces the uncertainty by 15%. The updated world average impacts the global precision electroweak fits, reducing the preferred $M_H$ by 6 GeV to $76^{+33}_{-24}$ GeV. The 95% CL upper limit on $M_H$ is $144_{-24}^{+182}$ GeV with (out) the LEP II direct limit included. Figure 2 shows the $M_T$ fits for $\Gamma_W$ in the muon and electron decay channels. The results are combined to give the final result $\Gamma_W = 2032 \pm 71$ MeV, the world's most precise single measurement, which is in good agreement with the SM prediction. This result reduces the world average central value by 44 MeV and uncertainty by 22%.
Figure 1: Transverse mass fits for \( M_W \) in \( W \rightarrow \mu\nu \) (left) and \( W \rightarrow e\nu \) (right) events. The fit is performed in the region 65–90 GeV.

Figure 2: Transverse mass fits for \( \Gamma_W \) in \( W \rightarrow \mu\nu \) (left) and \( W \rightarrow e\nu \) (right) events. The fit is performed in the region 90–200 GeV.

References

TOP PAIR PRODUCTION CROSS-SECTION AT THE TEVATRON

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An overview of latest top quark pair production cross-sections measured at the Tevatron is given. These measurements have been carried out in the dilepton, lepton+jets and all-jets channels with an integrated luminosity of about 1 fb$^{-1}$. The measurements are consistent with NNLO calculations.

Since the top quark discovery in 1995 by the CDF and DØ collaborations$^1$, top pair production cross-sections are one of the basic measurements to be carried out on each new data sample. During Tevatron Run I (1992-1996) an integrated luminosity of about 100 pb$^{-1}$ at a center of mass energy of $\sqrt{s} = 1.8$ TeV allowed to measure top pair production cross-sections of $6.5^{+1.7}_{-1.4}$ pb and $5.7 \pm 1.6$ pb by the CDF and DØ collaborations respectively. The Tevatron Run II started in 2001 and until spring 2006 about 1 fb$^{-1}$ of $pp$ collisions with $\sqrt{s} = 1.96$ TeV have been produced and analyzed since. At this energy an increase of about 30% in the cross-section is expected. The most recent NNLO calculations predict a cross-section of $6.7^{+0.7}_{-0.9}$ pb$^2$ or $6.8 \pm 0.6$ pb$^3$ for a top mass, $m_{t\bar{t}} = 175$ GeV.

In the standard model (SM) $|V_{tb}| \sim 1$ leads to a branching fraction $t \rightarrow Wb$ close to 100%. With a lifetime of $\tau \sim 10^{-25}$ s top quarks decay before hadronization. Their decay channels are classified according to the decay of the $W$ bosons produced. The dilepton channel accounts for about 6% of all decays, taking into account decays into $ee, \mu\mu$ and $e\mu$ and including leptonic $\tau$ decays. The $l+\text{jets}$ channel represents about 34% of the cross-section with the leptonic $\tau$ decays included in the $e+\text{jets}$ and $\mu+\text{jets}$ channels. Decays into all jets occur in 46% of the events, the remaining 14% correspond to signatures with hadronic $\tau$ decays.

Cross-section measurements are an important test of perturbative QCD at high $p_T$ as non-SM top production, for example resonant top production, may lead to higher a cross-section than expected. It is important to verify the consistency of different decay channels, as some non-SM
models for example $t \rightarrow H^+$ or $t \rightarrow \bar{t}$, modify the contributions in different decay channels.
Non-$W$ top decays are probed from the comparison of the dilepton and $l+$jets measurements. Top quark event selections using $b$-jet tagging assume a branching ratio $BR(t \rightarrow Wb) = 1$. Their consistency with kinematic methods, free of this assumption, is an important check of the SM prediction.

A cross-section is in most cases obtained from a counting experiment: $\sigma(pp \rightarrow t\bar{t}) = (N_{\text{obs}} - N_{\text{bkgd}})/A_{\text{tot}}L$. $N_{\text{bkgd}}$, the number of background events, estimated from Monte Carlo simulations and/or data samples, is subtracted from $N_{\text{obs}}$, the number of observed events meeting the selection criteria of a top-event signature. This difference is normalized by the integrated luminosity $L$ and the total acceptance $A_{\text{tot}}$. $A_{\text{tot}}$ includes the geometric acceptance as well as trigger efficiency and event selection efficiency and is slightly dependent on $m_{\text{top}}$. In all the Monte Carlo simulations $m_{\text{top}} = 175$ GeV has been used.

1 The dilepton channels

The signature of top dilepton events is two high $p_T$, opposite sign leptons ($p_T > 15$ GeV), some missing transverse energy ($E_T > 35$ GeV) and two or more, high $p_T$ jets ($p_T > 20$ GeV). Physics background is due dominantly to $Z/\gamma^*+\text{jets}$ events, $WW/ZZ$ events, and estimated from Monte Carlo simulations. Instrumental backgrounds occur due to fake isolated leptons, either as a mis-identified $e$ or a $\mu$ in a non-reconstructed $b$-jet, as well as $E_T$ from detector resolution, fake jets or noise in the calorimeter. These background are estimated from data.

D0 measured the cross-section in the $ee$, $\mu\mu$ and $e\mu$ channels. Requiring 2 leptons and 2 jets in the event selection yields to an acceptance of 8% and 5% in the $ee$ and $\mu\mu$ channels, and 12% in the $e\mu$ channel. The acceptance in $e\mu$ channel could be further improved by an additional 3% taking into account the events with only 1 jet in the final state. In total 73 events are observed for 51 expected signal events and 24 expected background events. Details for each channel are given in Table 1. The combined result from the three measurements is $\sigma_{t\bar{t}} = 6.8^{+1.2}_{-1.1}(\text{stat})^{+0.9}_{-0.8}(\text{syst}) \pm 0.4(\text{lumi})$ pb, with the main systematic errors being the lepton identification efficiency and the jet energy calibration.

An alternative method used by CDF loosens the lepton identification criteria for the second lepton, by requiring only an isolated track. The error on this measurement has been improved through an increase of the acceptance reaching 14%, even though the signal/background (S/B) ratio is reduced to 1.3. The number of expected and observed events are also given in table 1, leading to a cross-section of $\sigma_{t\bar{t}} = 9.0 \pm 1.3(\text{stat}) \pm 0.5(\text{syst}) \pm 0.5(\text{lumi})$ pb.

D0 carried out an exclusive lepton+track analysis on a data sample with 360 pb$^{-1}$, explicitly vetoing a fully reconstructed second lepton to allow for a combination with the dilepton measurements and using $b$-jet tagging to improve the purity of the sample. An update of this measurement with 1 fb$^{-1}$ is in progress.
2 The lepton+jets channel

For top decays into \(l^+\) jets, the signature is a high \(p_T\) lepton, large \(E_T\) and 4 or more high \(p_T\) jets. Dominating physics background is due to \(W+jets\) events. Instrumental background is due to fake isolated leptons in multijet events. To separate the signal and background in the \(l^+\) jets channels either the kinematic properties of the events are used or \(b\)-jet tagging is required.

For the first type of analysis, DØ constructs a likelihood discriminant based on six kinematic variables without a \(b\)-jet tagging requirement\(^6\). Its output is shown in figure 1(left) for the combined \(e^+\) jets and \(\mu^+\) jets sample. For an integrated luminosity of \(\mathcal{L} = 0.91\text{ fb}^{-1}\), 124 \((100)\) \(t\bar{t}\) events are expected in the \(e^+\) jets (\(\mu^+\) jets) channel for 168 \((235)\) \(W+jets\) events and 62 \((27)\) multijet events, leading to a combined cross-section of \(\sigma_{t\bar{t}} = 6.8^{+1.0}_{-0.8}\text{ (stat)+0.7\text{ (syst)+0.4\text{ (lumi)}}}\) pb. The dominant systematic errors come from the background model, lepton identification and jet energy scale.

The complementary method\(^7\) uses a new Neural Network (NN) based \(b\)-jet tagging algorithm. The chosen operating point has a \(b\)-jet tagging efficiency of 55\% and a fake tag rate of 1\%. For the same fake tag rate this represents a 15\% increase in efficiency with respect to the previous \(b\)-jet tagging algorithm. The cross-section result of \(\sigma_{t\bar{t}} = 8.3^{+0.6}_{-0.5}\text{ (stat)+0.9}_{-1.0}\text{ (syst)+0.5\text{ (lumi)}}\) pb is a combination from the NN-output in 8 different channels, considering \(e\) and \(\mu\) final states, 3-jet events, or 4 and more jet events, 1-tag or 2-tags separately. The jet-multiplicities for the combined 1-tag and 2-tags samples are shown in figure 1(middle and right respectively). Within the errors, the results are consistent between the kinematic and the \(b\)-tagging analyses.

CDF results in the \(l^+\) jets channels have been presented at the previous Moriond QCD conference\(^8\) on an integrated luminosity of \(\mathcal{L} = 0.7\text{ fb}^{-1}\). For a method using only kinematic variables the result is \(\sigma_{t\bar{t}} = 6.0 \pm 0.6\text{ (stat)+0.9\text{ (syst)+0.3\text{ (lumi)}}}\) pb. The analysis using \(b\)-jet tagging yields to a cross-section of \(\sigma_{t\bar{t}} = 8.2 \pm 0.6\text{ (stat)+0.9\text{ (syst)+0.5\text{ (lumi)}}}\) pb.

3 All jets channel

Even though this channel has the highest branching ratio, it is largely dominated by multijet background. The preselection of \(t\bar{t} \rightarrow jets\) in the inclusive CDF multijet sample\(^9\) requires 6 to 8 jets with \(p_T > 15\text{ GeV}\) separated by \(\Delta R > 0.5\). This preselection has a S/B ratio of 1/1300. To reduce the background a NN is used with 11 kinematic input variables. A further improvement in the S/B ratio is obtained by requiring a secondary vertex tag. The cross-section is then measured from the number of observed tags with an expectation of \(n_{\text{tag}} = 0.95 \pm 0.07\) per top event determined from Monte Carlo. In total 387 signal and 846 background events have been found for \(\mathcal{L} = 1.02\text{ fb}^{-1}\). The cross-section obtained is \(\sigma_{t\bar{t}} = 8.3 \pm 1.0\text{ (stat)+2.0}_{-1.5}\text{ (syst)+0.5\text{ (lumi)}}\) pb with the error on the jet energy calibration being the largest systematic contribution.
Figure 2: Top quark pair production cross-section measurements: summary of the current CDF (left) results and DØ results (middle). The right figure shows the $m_{top}$ dependence of the DØ dilepton measurement.

4 Summary

Summaries of the CDF and DØ top quark pair production cross-section measurements are shown in figure 2 (left and middle respectively). All the measurements are in good agreement with the NNLO-predictions. Results with an integrated luminosity of about 1 fb$^{-1}$ are highlighted. With this luminosity the errors have been sizably reduced and reach in some of the decay channels about 15%. From a combination of all results an experimental error at the order of the theoretical error can be expected.

During the conference the question was raised if $m_{top}$ could be determined from the cross-section measurements. A first answer concerning the precision that could be achieved can be interfered from figure 2 (right) showing the dependence of the DØ dilepton cross-section as a function of $m_{top}$. The shaded band shows the assumption of the total error for a combined result of the currently measured current cross-sections to be of the size of the theoretical error and with the same $m_{top}$ dependence than the di-lepton cross-section. A determination of $m_{top}$ using the current production cross-section would lead to an error on $m_{top}$ of about ±5 GeV. With the full Run II statistics an experimental error half this size looks a reasonable guess.

Acknowledgments

I would like to thank the CDF and DØ collaborations for presenting these results, in particular the top conveners Kirsten Tollefson, Robin Erbacher, Elizabeth Shabalina, Ulrich Heintz, Michele Weber for their help in preparing the talk, rehearsals and very useful comments, Kevin Lannon for a last minute plot and the organizers for another great Moriond!

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Top Quark Mass Measurements at the Tevatron and the Standard Model Fits

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New measurements of the top quark mass from the Tevatron are presented. Combined with previous results, they yield a preliminary new world average of $m_{\text{top}} = 170.9 \pm 1.1(\text{stat}) \pm 1.5(\text{syst}) \text{GeV}/c^2$ and impose new constraints on the mass of the Higgs boson.

1 Introduction

The huge interest in a precise measurement of the top quark mass ($m_{\text{top}}$) is primarily motivated by its role in constraining the mass of the Higgs boson ($m_{\text{Higgs}}$). To see this, let us begin by looking at the mass of the $W$ boson ($m_W$) in the Standard Model which, when one-loop radiative corrections are included, can be related to well known electroweak quantities through the following expression:

$$m_W^2 = \frac{\frac{3G_F}{4\pi^2}}{\sin^2 \theta_W (1 + \Delta r)}.$$  

(1)

The radiative corrections contained in $\Delta r$ receive contributions from the top quark:

$$(\Delta r)_{\text{top}} \approx \frac{3G_F m_{\text{top}}^2}{8\sqrt{2}\pi^2} \frac{1}{\tan^2 \theta_W}$$  

(2)

and the Higgs boson:

$$(\Delta r)_{\text{Higgs}} \approx \frac{11G_F m_Z^2 \cos^2 \theta_W}{24\sqrt{2}\pi^2} \ln \frac{m_{\text{Higgs}}^2}{m_Z^2}$$  

(3)

where $m_Z$ is the mass of the $Z$ boson. From these expressions, we see that $m_{\text{top}}$ enters quadratically while $m_{\text{Higgs}}$ enters logarithmically. A precise knowledge of both $m_W$ and $m_{\text{top}}$ in combination with existing electroweak data is therefore necessary to impose useful constraints on $m_{\text{Higgs}}$. Such constraints, in turn, are of tremendous value in the ongoing search for the Higgs.

In this talk, we present the latest top quark mass measurements from the CDF and DØ collaborations based on up to 1 fb$^{-1}$ of Run II data collected at Fermilab’s Tevatron. These results are combined with previous ones to give a new preliminary world average for $m_{\text{top}}$ which, in turn, yields new constraints on the Higgs mass.
2 Measurement Channels and Experimental Challenge

Now that we understand the motivation behind a precise determination of the top mass, let us look at the top quark decay channels in which these measurements are performed and the experimental challenges they pose.

In the all jets channel, both $W$ bosons from the $t\bar{t}$ pair decay hadronically into jets for a total of 6 jets in the event. This channel has the advantage of having the largest branching ratio of 44%. It suffers, however, from large background levels from QCD multijet events. On the other hand, it benefits from the presence of the hadronically decaying $W$ bosons whose well known masses can be exploited to perform an in-situ calibration of the jet energies, reducing the effect of the systematic uncertainty in the overall jet energy scale. In the dilepton channel, both $W$ bosons decay leptonically. It has the advantage of having the lowest background levels coming from Drell-Yan processes associated with jets, diboson production with associated jets, and $W + j$ events with one jet faking an electron. Unfortunately, it also has the lowest branching ratio of 5%. In the lepton+jets ($\ell$+jets) channel, one of the two $W$ bosons from the $t\bar{t}$ pair decays hadronically while the other one decays leptonically. This channel maintains a good balance between a reasonable branching ratio of 29% and moderate background levels from $W + j$ and QCD multijet events. Like the all jets channel, it can benefit from an in-situ jet energy calibration using the $m_{\ell\nu}$ constraint. It has traditionally yielded the most precise $m_{\text{top}}$ measurements.

To appreciate the challenge involved in measuring the top mass at the Tevatron, let us now take the $\ell$+jets channel as an example. In this case, what our reconstruction programs give us from the detector are several jets, a high $p_T$ lepton, substantial missing transverse energy, and an interaction vertex. Since we don't really know how to associate jets with partons in general, all jet permutations need to be considered in a straightforward reconstruction of the top mass. Furthermore, unlike long lived particles, there are no detached vertices associated with the top quark itself that can be used to separate the signal from the background events. This means that, even with $b$-tagging, there are no sharp and clean mass peaks from which the top mass can be determined directly. Fortunately, despite these challenges, sophisticated measurement techniques have been developed that make a precise measurement of the top mass possible.

3 Top Quark Mass Measurement Techniques

In this section we describe the three major techniques used in measuring the top quark mass. All the measurements presented here use one or some combination of these techniques.

The template method is the oldest of the three techniques and has been used for most of the earliest mass measurements. In this technique, one begins by identifying a variable sensitive to the top mass, an obvious choice of which would be the kinematically reconstructed value of the mass itself. Distributions of the chosen variable are then plotted separately for several samples of fully simulated Monte Carlo (MC) events differing only in the value of the top mass used to generate the signal events. Each of these distributions is called a template and is associated with a particular value of the input mass. The top mass is then extracted from the data sample by comparing the data distribution directly with each MC template to find the best fit value based on some measure of the goodness of fit. More recent applications of this technique parameterize the templates in terms of a probability density function which is used to construct likelihoods from which the top mass is extracted.

$a^{0}$ pioneered the application of the matrix element (ME) method to top quark mass measurements in the Run I data from the $\ell + j$ channel. It is based on calculating the probability for observing each event which includes contributions from both signal and background sources. The signal probability is calculated as a function of the assumed top mass, resulting in a prob-
ability distribution for each event. The probability is taken to be the differential cross section for the process in question. The calculated probability distributions for every event in the data sample are combined to construct a joint likelihood from which the top mass is determined and its uncertainty estimated. The ME method makes use of as many measured variables as possible to completely specify an event, thereby allowing maximum discrimination between signal and background events. Within each event, all possible jet permutations are combined in a natural way based on their relative probabilities. Furthermore, the use of transfer functions allows a probabilistic treatment of the mapping between parton and jet energies where the full spectrum of parton energies contributing to the observed jet energy is taken into account.

The ideogram method, like the ME method, calculates an event-by-event likelihood. This technique makes use of a constrained kinematic fit to reconstruct the top mass. Using a simple parameterization, the probability for observing the reconstructed mass is then calculated as a function of the true value with the measurement resolution taken into account. This technique, which was also pioneered by DØ\textsuperscript{2}, aims to achieve statistical uncertainties comparable to those of the ME method without requiring as many computational resources.

4 New Results from the Tevatron

DØ has measured the top quark mass in the $\ell+$jets channel using the ME method described in the previous section\textsuperscript{3}. This measurement takes advantage of the $m_W$ constraint to perform an in-situ calibration of the jet energies. This is done by introducing a global scale factor, $JES$, that is applied to the energies of all the jets. A fit is then performed that maximizes the likelihood simultaneously in $m_{\text{top}}$, $JES$, and the signal fraction $C_s$. The 2D likelihood fits in $m_{\text{top}}$ and $JES$ are shown separately for the electron and muon channels in Figure 1. The combined result for both channels is $170.5 \pm 2.4(\text{stat} + \text{JES}) \pm 1.2(\text{syst})\text{GeV}/c^2$ for 0.9 fb\textsuperscript{-1} of data. Dominant systematic uncertainties are in the modeling of initial and final state radiations and $b$-fragmentation. This is the best DØ measurement of the top quark mass to date.

CDF has also measured the top quark mass in the $\ell+$jets channel using the ME method. Like the DØ result, this measurement employs an in-situ jet energy calibration through the inclusion of a global $JES$ parameter in the likelihood fit. The left plot in Figure 2 shows the 2D likelihood fit to the data for both electron and muon channels in $JES$ and $m_{\text{top}}$. The right plot in Figure 2 shows the expected error distribution from MC ensemble tests with the arrow indicating the measurement uncertainty. The measured result for 0.94 fb\textsuperscript{-1} of data is $170.9 \pm 2.2(\text{stat} + \text{JES}) \pm 1.4(\text{syst})\text{GeV}/c^2$. The largest systematic uncertainty is in the modeling of initial and final state radiations. This is currently the most precise CDF measurement of the
top quark mass.

DØ has a measurement of the top quark mass in the \(\ell+\text{jets}\) channel using the ideogram method\(^2\). Like the two results above, it employs an in-situ jet energy calibration. The 2D likelihood as a function of \(JES\) and \(m_{\text{top}}\) is shown on the left in Figure 3 with the gray line indicating the fitted value of \(JES\) as a function of \(m_{\text{top}}\). The right plot in Figure 3 shows the 1D likelihood as a function of \(m_{\text{top}}\) along the gray line in the left plot. The result for 0.4 fb\(^{-1}\) of data is \(173.7 \pm 4.4\) (stat + JES)\(^{+2.1}_{-2.0}\) (syst) GeV/c\(^2\). Dominant systematic uncertainties are in the modeling of \(b\)-fragmentation and in the \(b/\text{light jet}\) energy scale ratio.

CDF has applied the ME method to a measurement of the quark top mass in the dilepton channel\(^5\). A plot of the probability as a function of \(m_{\text{top}}\) is shown on the left in Figure 4 and the expected error distribution on the right with the arrow indicating the measurement uncertainty. The result for 1 fb\(^{-1}\) of data is \(164.5 \pm 3.9\) (stat) \(\pm 3.9\) (syst) GeV/c\(^2\). The systematic error is dominated by the uncertainty in the jet energy scale.

DØ has measured the top quark mass in the dilepton channel using a template method that assigns a weight to each neutrino solution based on the agreement between the calculated transverse momentum of the neutrinos and the observed missing transverse energy\(^6\). The result for 1 fb\(^{-1}\) is \(172.5 \pm 5.8\) (stat) \(\pm 5.5\) (syst) GeV/c\(^2\). The dominant source of the systematic error is the jet energy scale uncertainty.

CDF has measured the top quark mass in the all jets channel using a combination of template and ME methods\(^7\). Instead of using the ME method directly to measure the top mass, the value
determined from the method is used to construct the MC templates. Probabilities calculated from the ME are also used in the event selection process to identify events with high signal probability. This result also uses the $m_W$ constraint to perform an in-situ jet energy calibration. The left plot in Figure 5 shows a fit of the data distribution to the MC templates for events with two $b$-tagged jets. Contours of $JES$ and $m_{top}$ in data are shown on the right in Figure 5. The result for 1 fb$^{-1}$ is $171.1 \pm 3.7(\text{stat} + \text{JES}) \pm 2.1(\text{syst})$ GeV/c$^2$. The largest systematic uncertainties are in the simulation of fragmentation and showering and of final state radiation.

5 New World Average and Standard Model Fits

From above, the best result of each experiment in each channel is combined with previous results yielding a new preliminary world average of $m_{top} = 170.9 \pm 1.1(\text{stat}) \pm 1.5(\text{syst})$ GeV/c$^2$ shown on the left in Figure 6. The ME $\ell$+jets results from DØ and CDF carry the largest weights in this average of 40% and 39%, respectively. This value is 0.5 GeV/c$^2$ lower than the previous world average. With this new preliminary result, the top quark mass is now known to a total uncertainty of 1.8 GeV/c$^2$ corresponding to a relative precision of 1.1%.

This new top quark mass is also combined with other precision electroweak results in Standard Model fits performed by the LEP Electroweak Working Group. The right plot in Figure 6 shows the $\Delta \chi^2$ curve resulting from these fits giving $m_{Higgs} = 76.1^{+33}_{-24}$ GeV/c$^2$ at the minimum and a 95% confidence level upper limit of 144 GeV/c$^2$ which increases to 182 GeV/c$^2$ when the
LEP-2 direct search limit of 114 GeV/c² indicated by the yellow band is included.

6 Summary and Conclusions

A precise determination of the top quark mass is crucial for constraining the mass of the Higgs boson. Despite the great challenges involved, precise measurements are possible through the use of sophisticated measurement techniques. This talk presented new results based on up to 1 fb⁻¹ of data collected by CDF and DØ. Although these results are still dominated by the ℓ+jets channel, the other two show promise and we hope to see more competitive results from them in the future. Combining the new results with previous ones has yielded a new preliminary world average top quark mass with a total uncertainty of 1.8 GeV/c² and imposed new constraints on $m_{Higgs}$. As more data become available at the Tevatron, we can expect statistical uncertainties < 1 GeV/c² by the end of the Tevatron run at which point the total uncertainties will become dominated by the systematic uncertainties.

References

MEASUREMENTS OF TOP PROPERTIES AT THE TEVATRON

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The large data samples of thousands of top events collected at the Tevatron experiments CDF and DØ allow for a variety of measurements to analyze the properties of the top quark. Guided by the question “Is the top quark observed at the Tevatron really the top quark of the standard model?,” we present Tevatron analyses studying the top production mechanism including resonant $t\bar{t}$ production; the $V-A$ structure of the $t\rightarrow Wb$ decay vertex, the charge of the top quark, and single-top production via flavor-changing neutral currents.

Keywords: Hadron Collider, Tevatron, Heavy Quark Production, Top Quark Properties

1 Introduction

At the Tevatron collider at Fermi National Accelerator Laboratory, protons and antiprotons are collided at a center-of-mass energy of 1.96 TeV. The Tevatron provides the highest energies currently available at a hadron collider and is the only collider with sufficient energy to produce top quarks. The two multi-purpose experiments at the Tevatron, CDF $^{1}$ and DØ $^{2}$, have studied production and decays of top quarks in great detail. The focus of this article will be on measurements of top properties at CDF and DØ (other than mass and cross section). The event selection for top properties measurements is built on the experience gained in the mass and cross section analyses in the “lepton+jets” channel ($t\bar{t}\rightarrow WbWb\rightarrow \ell\nu b\ell\nu b$) and the “dilepton” channel ($t\bar{t}\rightarrow WbWb\rightarrow \ell\nu b\ell\nu b$). Also the background composition, with $W$ production in association with jets as the dominant source, and systematic uncertainties, mainly coming from determining the jet energy scale, are similar to those in the top mass and cross section analyses.

2 Top Pair Production

2.1 Fraction of $t\bar{t}$ Pairs Produced via Gluon-Gluon Fusion

In the standard model (SM), the production of $t\bar{t}$ pairs at the Tevatron is dominated by the process of $q\bar{q}$ annihilation. The contribution of the $gg$ fusion channel to the $t\bar{t}$ production cross section amounts to $15 \pm 5\%$, where the large uncertainty is due to poor knowledge of the gluon parton distribution function inside the proton $^{3,4}$. The CDF collaboration has developed two complementary methods to measure the fraction of $t\bar{t}$ pairs produced via $gg$ fusion. With datasets of up to 1 fb$^{-1}$ of integrated luminosity, both methods are dominated by statistical uncertainties and are therefore expected to improve with more data.

One method $^{5}$ utilizes an artificial neural network (NN) to distinguish the processes $q\bar{q}\rightarrow t\bar{t}$, $gg\rightarrow t\bar{t}$, and $q\bar{q}/gg\rightarrow W$+jets. The NN input comprises the velocity and angle of the top quark, and the angles of the three decay products in the off-diagonal spin basis. From the resulting NN discriminant, an upper limit on the $gg$ fraction is derived via a Feldman-Cousins method that includes systematic uncertainties. As shown in Fig. 1a, the measurement yields a limit of $\sigma(gg\rightarrow t\bar{t})/\sigma(pp\rightarrow t\bar{t}) < 0.51$ at 95% C.L.
Figure 1: Measurements of the $gg$ fraction in $t\bar{t}$ production. (a) Feldman-Cousins band obtained for neural network discriminant method. (b) Distribution of low-$p_T$ tracks fitted by gluon-rich and no-gluon templates in $W+4$-jet data.

The other method\(^6\) is based on the observation that the number of tracks with small transverse momenta $p_T$ in the range of $0.3-2.9\text{GeV}/c$ is strongly correlated with the number of gluons in the event. The correlation is calibrated using gluon-rich and gluon-free control samples in the data, and then extrapolated to the top-rich sample of $W+4$ jets. Gluon-rich and no-gluon templates are fitted to the distribution of low-$p_T$ tracks in the data, as depicted in Fig. 1b, to extract a $gg$ fraction of $\frac{\sigma(gg \rightarrow t\bar{t})}{\sigma(p\bar{p} \rightarrow t\bar{t})} = 0.01 \pm 0.16(\text{stat.}) \pm 0.07(\text{syst.})$.

2.2 Resonant Production of $t\bar{t}$ Pairs

Several extensions of the SM predict $t\bar{t}$ production from the decay of heavy particles. In one particular model\(^7\), the heavy particle is assumed to be a narrow $Z'$ resonance that couples strongly only to third generation quarks and does not couple to leptons ("leptophobic $Z'$").

The CDF collaboration has studied the production of a $Z'$-like heavy $t\bar{t}$ resonance\(^6\). The $t\bar{t}$ invariant mass $M_{t\bar{t}}$ is reconstructed with the help of a kinematic fitter developed for top mass measurements. The $M_{t\bar{t}}$ distribution in the data is compared to the expectation for SM $t\bar{t}$ production and non-t$^\ell$ background. The background includes $W+$jets events, QCD multijet events, and diboson production ($WW$, $WZ$, $ZZ$), and was estimated with a method that combines input from Monte Carlo simulations and data control samples. The data are compatible with SM $t\bar{t}$ production, so that a limit on the $Z'$ production cross section can be derived, see Fig. 2a. From the comparison with the cross section prediction for leptophobic $Z'$, CDF obtains a limit of $M_{Z'} > 725 \text{ GeV}/c^2$ at 95% C.L.

3 Helicity of $W$ Bosons from Top Decays

Due to its small lifetime of less than $10^{-24}$ s, the top quark decays before it hadronizes. As a consequence, the complete spin information is transferred to its daughter particles. In the SM, top quarks decay to a $W$ boson and a $b$ quark almost exclusively. The spin-1 $W$ boson has three possible helicity states: longitudinal, left-handed, and right-handed. The $V-A$ structure of the $tWb$ vertex in the SM does not allow for a right-handed state, so that any sizable admixture of right-handed $W$'s would be a sign of new physics. The longitudinal fraction of the $W$'s helicity in the SM is determined by the Yukawa coupling of the top quark and amounts to $f^0 \approx 0.7$.

Both CDF and DØ have measured the right-handed helicity fraction $f^+$ of $W$ bosons from top decays\(^9,10,11,12\). The analyses use different techniques to reconstruct the observable $\cos \theta^*$, the angle between the top boost direction and the charged lepton in the $W$ rest frame, or the
correlated quantity $M_{lb}$, the invariant mass of the charged lepton and the $b$ jet. Due to the limited size of the data samples, most analyses\textsuperscript{9,10,11} check the consistency of $f^0$ with the SM, assuming $f^+=0$, and fix the value to $f^0 = 0.7$ before $f^+$ is measured. With 1 fb$^{-1}$ of data, a simultaneous fit to $f^0$ and $f^+$ becomes feasible\textsuperscript{12}. As shown in Fig. 2b, all measurements are compatible with the SM prediction of $f^+=0$.

4 Top Charge

In the SM, the top quark and the bottom quark form a left-handed isospin doublet with charges $(2/3 \, e, -1/3 \, e)_L$, where $e$ is the elementary charge. Fits to electroweak precision data can be improved using a theoretical model with a top mass of 270 GeV/$c^2$ and an exotic right-handed quark doublet with charges $(-1/3 \, e, -4/3 \, e)_R$ that mixes with right-handed $b$ quarks\textsuperscript{13}. In this model, the exotic replacement for the top quark has a charge of $-4/3 \, e$.

The CDF collaboration has performed a measurement to test if the top charge is $2/3 \, e$ or $-4/3 \, e$\textsuperscript{14}. The measurement consists of three steps. The $W$ charge is obtained via the charge of the lepton in the decay $W \to \ell \nu$. To reconstruct the top from the decay $t \to Wb$, the $W$ is paired with a $b$ jet, using a kinematic fitter (lepton+jets channel) or a cut on $M_{lb}$ (dilepton channel). Finally, the observable “jet charge,” a weighted sum of the charge of all tracks that form a jet, is employed to measure the flavor of the $b$ jet. The product of $W$ charge and jet charge is used to distinguish the exotic model from the SM, see Fig. 3a.

CDF has tested the consistency of the data with the SM and the exotic model with a hypothesis test, with the null hypothesis that the SM is correct. If the exotic model were correct, 81% of all measurements would return $p$-values below 0.01 (probability to incorrectly reject the SM, chosen \textit{a priori}). The measured $p$-value of 0.35 shows that the data are consistent with the SM, and the exotic model is excluded at 81% C.L.

5 Single-Top Production via Flavor-Changing Neutral Currents

Flavor-changing neutral currents (FCNC) in the top quark sector are heavily suppressed in the SM. For example, the branching fraction $\mathcal{B}(t\to gc)$ is approximately $5 \times 10^{-12}$, far below the reach of present and future hadron collider experiments\textsuperscript{15}. Any FCNC signal at the Tevatron would be a sign of physics beyond the SM.
Figure 3: (a) Top charge measurement via jet charge. The data are compared to the predictions of the SM for $t\bar{t}$ signal and background. (b) Exclusion contours for the couplings $(\kappa/\Lambda)^2$ of the flavor-changing neutral current vertices $tgu$ and $tgc$.

The DØ collaboration has searched for FCNC in the production of single top quarks\textsuperscript{16}. In the presence of the processes $ug \to t$ and $cg \to t$, the single top production rate is enhanced. DØ has deployed a NN to discriminate the kinematics of the FCNC signal from $t\bar{t}$, $W+$jets, and QCD multijet backgrounds. The data agrees very well with the SM prediction, so that an upper limit on the FCNC couplings $tgu$ and $tgc$ could be derived. Fig. 3b shows the upper limit as a function of the strengths $(\kappa/\Lambda)^2$ of the $tgu$ and $tgc$ coupling, where $\kappa$ is the coupling parameter in the Lagrangian, and $\Lambda$ is a generic new physics scale. The DØ measurement improves upon previous measurements, prominently by the HERA experiments, by a factor of 3–11.

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TOP PAIR PRODUCTION AT THRESHOLD

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I give an overview of the third-order calculations for the heavy-quarkonium parameters in the nonrelativistic effective theory framework and its application to the phenomenology of top quark threshold production. The focus is on the ultrasoft contribution\(^1\).

1 Introduction

Theoretical study of nonrelativistic heavy quark-antiquark systems is among the earliest applications of perturbative quantum chromodynamics (QCD)\(^2\) and entirely relies on the first principles. The nonperturbative effects are well under control and the reliable theoretical predictions can be obtained within the perturbation theory. This makes the heavy quark-antiquark systems an ideal laboratory to determine fundamental parameters of QCD, such as the strong coupling constant \(\alpha_s\) and the heavy-quark mass \(m_Q\).

The binding energy of the heavy-quarkonium state and the value of its wave function at the origin are among the characteristics of the heavy quarkonium system that are of primary phenomenological interest. The former determines the mass of the bound state resonance, while the latter controls its production and annihilation rates. The heavy-quarkonium ground state energy has been computed through \(O(\alpha_s^2m_Q)\) including the third order correction to the Coulomb approximation in Refs.\(^3,4,8\) This result has been extended to the excited states in Refs.\(^9,10\) For the wave function at the origin however a complete result is only available through \(O(\alpha_s^2)\)\(^6,7,8,9,10,11,12\) The second order correction is huge and for a reliable perturbative prediction one has to go full control over the next order. In this order the complete result is available only for logarithmically enhanced terms which include the double logarithmic \(O(\alpha_s^3 \ln^2 \alpha_s)\) contribution\(^13\) and
the single logarithmic $O(\alpha_s^3 \ln \alpha_s)$ contribution \cite{14,15} (see also Refs. \cite{16,17,18}). The calculation of the most difficult non-logarithmic term has been started in Refs. \cite{19,20}, where the contribution to the wave function at the origin from the loop corrections to the Coulomb potential have been evaluated. In Ref. \cite{21} the contributions from the non-Coulomb potentials have been obtained. The last breakthrough is the calculation of the contribution due to the emission and absorption of an ultrasoft gluon by the quarkonium bound state, \cite{1} which completes the analysis of the nonrelativistic quarkonium bound-state dynamics in the third order. The full third-order correction to the wave function at the origin is now expressed in terms of a few yet unknown matching coefficients, which can be obtained by standard fixed-order loop calculations. In this paper I outline the effective theory approach to the theory of heavy quarkonium and present the result for the ultrasoft contribution to the top quark-antiquark resonance cross section. The ultrasoft correction is of special interest, because it constitutes a qualitatively new effect, which shows up for the first time in the third order. No other such effects are expected in higher orders of the perturbative expansion.

2 Nonrelativistic effective theory

Near the threshold, the heavy quarks are nonrelativistic, so that one may consider the quark velocity $v$ (or inverse quark mass) as a small parameter. An expansion in $v$ may be performed directly in the QCD Lagrangian by using the framework of effective field theory, \cite{22,23,24} or diagrammatically with the threshold expansion. \cite{25} The relevant momentum regions are the hard region (energy $k^0$ and momentum $k$ of order $m$), the soft region ($k^0, k \sim mv$), the potential region ($k^0 \sim mv^2, k \sim mv$), and the ultrasoft region ($k^0, k \sim mv^2$). By integrating out the hard modes of QCD, one arrives at the effective theory of nonrelativistic QCD (NRQCD). \cite{23} If one also integrates out the soft modes and the potential gluons, one obtains the effective theory of potential NRQCD (pNRQCD), which contains potential heavy quarks and ultrasoft gluons as dynamical particles. \cite{24} The propagation of the quark-antiquark pair in pNRQCD is described by the Green function of the Schrödinger equation

$$ (\mathcal{H} - E) G(r, r', E) = \delta^{(3)}(r - r'), $$

where $\mathcal{H}$ is the nonrelativistic effective Hamiltonian of the following form

$$ \mathcal{H} = -\frac{\partial^2}{m_q^2} + V_C(r) + \ldots, $$

with $r = |r|$. The ellipses stand for the higher order terms in $\alpha_s$ and $v$. For the color singlet state the leading order Coulomb potential is attractive, $V_C(r) = -C_F \alpha_s/r$, where $C_F = (N_c^2 - 1)/(2N_c)$, $N_c = 3$. As a consequence the color singlet Green function gets a contribution from an infinite number of approximately Coulombic bound states of the following form:

$$ G(r, r', E) = \sum_{n=1}^{\infty} \frac{\psi_n^*(r) \psi_n(r')}{E_n - E - i\epsilon} + \ldots, $$

where $E_n$ and $\psi_n$ are the energy and the wave function of a bound states, respectively, $n$ is the principal quantum number, the spin and orbital quantum numbers are suppressed and the ellipsis stands for the contribution of the spectral continuum. The leading order approximation of the quarkonium bound state is given by the Coulomb solution of Eq. (1), e.g. the leading order binding energy is $E_n^C = -C_F \alpha_s m_q/(2n)^2$. The corrections due to the high order terms in the nonrelativistic Hamiltonian can be systematically computed by means of the time ordered quantum mechanical perturbation theory. In addition, the Green function gets the correction due to the multipole interaction of the quark-antiquark pair to the ultrasoft gluons. The leading ultrasoft effect is due to the chromoelectric dipole interaction, which results in a NNNLO correction. \cite{26,27} The PNRQCD diagram representing this correction is shown in Fig. 1.
3 Ultrasoft contribution to top-quark production near threshold

For top quarks the nonperturbative effects are negligible and its decay width $\Gamma_t \approx 1.4$ GeV smears out the Coulomb resonances below the threshold. The NNLO analysis of the cross section shows that only the ground-state pole gives rise to a prominent resonance. Although the calculation of the normalized cross section $R = \sigma(e^+e^- \to t\bar{t}X)/\sigma(e^+e^- \to \mu^+\mu^-)$ requires the full Green function the height of the resonance can be estimated from the wave function at the origin of the would-be toponium ground state. In the leading-order approximation $R_1^{1,0} \approx 6\pi Ne^2/l^2|\psi_0^C(0)|^2/(m_t^2\Gamma_t)$, where $|\psi_0^C(0)|^2 = (m_t\alpha_s C_F)^3/(8\pi n^3)$ is the value of the Coulomb wave function at the origin. The ultrasoft correction to the wave function results in the following variation of the resonance cross section

$$\delta^{\alpha_s} R_1 = \alpha_s^3 \left\{ -18.71 \ln^2 \alpha_s + 52.03 \ln \alpha_s + 112.38 \right. \\
+ \left[ 23.52 \ln \alpha_s - 30.98 \right] \ln \frac{\mu}{m} - 6.55 \ln^2 \frac{\mu}{m} \right\} R_1^{1,0}. \quad (4)$$

The scale of the coupling $\alpha_s$ is most naturally of order of the inverse Bohr radius $m_t\alpha_s C_F$ in two of the three powers of the overall factor $\alpha_s^3$, and of order of the ultrasoft scale $m_t\alpha_s^2$ in the third. However, any other scale choice is formally equivalent at this order. In the following we evaluate $\alpha_s$ at $\mu_B = m_t C_F \alpha_s(\mu_B)$, wherever it appears. The scale $\mu$ in the $\ln(\mu/m_t)$ terms is related to scale-dependent potentials and hard matching coefficients. We vary $\mu/m_t$ between $\alpha_s C_F$ (corresponding to the scale $\mu_B$) and 1 (hard scale). Adopting $\alpha_s = 0.14$, which corresponds to $\mu_B \approx 32.5$ GeV, we obtain $\delta^{\alpha_s} R_1/R_1 \approx 0.31$ from the nonlogarithmic correction alone. Including an estimate of the logarithmic terms we find

$$\delta^{\alpha_s} R_1 \approx \left[ (-0.17) - (+0.13) \right] R_1^{1,0}. \quad (5)$$

It therefore appears that the large nonlogarithmic term leads to a large enhancement of the width. Whether or not perturbation theory is out of control (as may be suggested by the upper limit of the given range) can be decided only after combining all third-order terms.

4 Summary

The problem of evaluating the total $O(\alpha_s^3)$ corrections to the top quark threshold production is reduced to the fixed-order loop calculation of of a few yet unknown matching coefficients in dimensional regularization. The nonlogarithmic ultrasoft contribution is large and significantly increases the production rate. It might limit the accuracy of the perturbative analysis of the quarkonium even for top quarks. We should however emphasize that a definite conclusion can only be drawn once the full NNNLO result is available. In this respect the sizable negative third-order correction from the perturbation potentials should be mentioned.
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References

Higgs Amplitudes From Twistor Inspired Methods

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We illustrate the use of new on-shell methods, 4-dimensional unitarity cuts combined with
on-shell recursions relations, by computing the $A_4^{(1)}(\phi, 1^-, 2^-, 3^+, 4^+)$ amplitude in the large
top mass limit where the Higgs boson couples to gluons through an effective interaction.

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1 Introduction

The time for experiments at the LHC is approaching fast and it will be extremely important to
have accurate predictions of Standard Model processes if we hope to find signals of new physics.
Recent advances in on-shell techniques1,2,3 have made it possible to calculate compact analytic
expressions for multiparticle scattering amplitudes at one loop. Here we consider the application
of these so called “twistor” inspired methods to one loop amplitudes with a massive, colourless
scalar (the Higgs boson) two negative helicity gluons and two positive helicity gluons. This
builds on work for amplitudes with simpler helicity configurations4,5,6.

We consider the leading colour contribution to the one-loop amplitudes and, following3, we
split them into evaluating the “pure” 4-dimensional cut-constructible $C_4$ and rational $R_4$.

$$A_4^{(1)}(\phi, 1^-, 2^-, 3^+, 4^+) = C_4(\phi, 1^-, 2^-, 3^+, 4^+) + R_4(\phi, 1^-, 2^-, 3^+, 4^+). \quad (1)$$

The pure cut piece contains all (poly)logarithmic terms ($\log, \text{Li}_2, \pi^2$) which can be found by
computing the unitarity cuts in 4 dimensions7,8. The remaining rational terms can then be
evaluated using on-shell recursion relations.

2 The Model

In the Standard Model the Higgs boson couples to gluons through a fermion loop where the
dominant contribution is from the top quark. It is well known that for large $m_t$, the top quark
loop can be integrated out leading to the effective interaction,

$$L_{\text{int}}^\text{eff} = \frac{C}{2} H \text{tr} G_{\mu\nu} G^{\mu\nu}. \quad (2)$$

In the Standard Model, and to leading order in $\alpha_s$, the strength of the interaction is given by
$C = \frac{g^2}{64\pi^2} + O(\alpha_s^2)$, with $v = 246$ GeV. $C$ has been calculated up to order $O(\alpha_s^4)9,10$.

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*Analytic expressions for the $\phi$-MHV amplitude an arbitrary number of positive helicity gluons can be found
in4
The MHV structure of the Higgs-plus-gluons amplitudes is best elucidated \(^{10}\) by considering \(H\) to be the real part of a complex field \(\phi = \frac{1}{2}(H + iA)\), so that

\[
\mathcal{L}_{H,A}^\text{int} = \frac{C}{2} \left[ H \, \text{tr} \, G_{\mu\nu} G^{\mu\nu} + iA \, \text{tr} \, G_{\mu\nu} \, G^{\mu\nu} \right] = \frac{C}{2} \left[ \phi \, \text{tr} \, G_{SD \mu\nu} G^{SD}_{\mu\nu} + \phi^\dagger \, \text{tr} \, G_{ASD \mu\nu} G^{ASD}_{\mu\nu} \right]
\]

(3)

The amplitudes of the \(\phi\) and \(\phi^\dagger\) turn out to be much simpler than the corresponding \(H\) and \(A\) fields and so we proceed by calculating helicity amplitudes for gluons coupling to the \(\phi\) and then construct the \(\phi^\dagger\) amplitudes using parity symmetry. The full Higgs amplitudes are then made from the sum of \(\phi\) and \(\phi^\dagger\) amplitudes.

3 Cut Constructible Contributions

The unitarity method relies on sewing together tree-level amplitudes with on-shell propagators. Here we use the method of Brandhuber, Spence and Travaglini\(^ {11}\) which uses the off-shell continuation used for the tree-level MHV rules\(^ {12}\) to sew together tree MHV amplitudes. The subsequent integration over the continuation parameter, \(z\), reconstructs only the specific parts of the integral functions which have cuts in the considered channel. The cut integrals that are encountered have been considered previously by Van Neerven\(^ {13}\).

![Figure 1: The topologies for the cut-constructable part of the four gluon \(\phi\)-MHV amplitude.](image)

The tree-level QCD amplitudes have been known for some time\(^ {14}\) and the relevant \(n\)-point tree-level \(\phi\)-amplitudes have been recently computed using the CSW method\(^ {10,15}\). The topologies obtained by sewing together all possible configurations of joining a tree-level \(\phi\)-MHV amplitude with a pure QCD amplitude are shown in figure 1 where both gluons and fermions are allowed to circulate in the loop\(^ 4\). The final result is given by\(^ 4\):

\[
C_4(\phi, 1^-, 2^-, 3^+, 4^+) = \alpha_A A_4^{(0)}(\phi, 1^-, 2^-, 3^+, 4^+) \left[ U_4 + \left( \frac{N_P}{3} \frac{1432}{(12)^3} L_3(s_{341}, s_{41}) + \frac{N_P}{3} \frac{2341}{(21)^3} L_3(s_{234}, s_{23}) \right) \right.
\]

\[
- \frac{N_P}{2} \frac{1432}{(12)^2} L_2(s_{341}, s_{41}) - \frac{N_P}{2} \frac{2341}{(21)^2} L_2(s_{234}, s_{23})
\]

\[
+ \frac{N_P}{6} \frac{1432}{(12)} L_1(s_{341}, s_{41}) + \frac{N_P}{6} \frac{2341}{(21)} L_1(s_{234}, s_{23})
\]

\[
+ \frac{\beta_0}{N} \frac{1432}{(12)} L_1(s_{341}, s_{41}) + \frac{\beta_0}{N} \frac{2341}{(21)} L_1(s_{234}, s_{23}) \right]
\]

(4)
where for convenience, we have introduced
\[
\beta_0 = \frac{11N - 2N_F}{3}, \quad N_F = 2 \left(1 - \frac{N_F}{N}\right), \quad L_k(s,t) = \frac{\log(s/t)}{(s-t)^k}
\]
(5)
and
\[
U_4 = \frac{1}{2} \sum_{i=1}^{4} F_{i}^{3}m_{3}(s_{i+1},s_{i+2}) - F_{i}^{3}m_{3}(s_{i+1},s_{i+2}) - \frac{1}{2} \sum_{i=1}^{4} F_{i}^{4}m_{4}(s_{i+1},s_{i+2},s_{i+3},s_{i+4})
\]
(6)
The one-mass triangle $F_{i}^{3}m_{3}$ and box functions $F_{i}^{4}m_{4}$ can be found in $^7$.

4 Rational Contributions

We calculate the rational part using the unitarity bootstrap proposed by Bern, Dixon and Kosower$^3$ which generalised the tree level recursion of Britto, Cachazo and Feng$^6$. The method relies on simple complex analysis and the factorisation properties of one-loop amplitudes$^{17}$. In order to use this method it is very important to first remove all spurious singularities from the pure cut terms. Once this is achieved the remaining rational terms can be calculated using a recursion relation.

The spurious poles in eq. (4) appear in the functions $L_2$ and $L_3$, which can be removed by replacing these functions with new functions $\tilde{L}_2$ and $\tilde{L}_3$
\[
L_i(s,t) = \tilde{L}_i(s,t) + \frac{1}{2(s-t)^{i-1}} \left( \frac{1}{t} + \frac{1}{s} \right), \quad i = 2, 3.
\]
(7)
The additional rational terms must then be subtracted off again such that $C_4 = \tilde{C}_4 + CR_4$ with $CR_4$ given by:
\[
CR_4(\phi,1^-,2^-,3^+,4^+) = C_4[L_1,L_2,L_3] - C_4[\tilde{L}_1,\tilde{L}_2,\tilde{L}_3]
\]
(8)
The direct recursive terms are calculated by making a shift into complex momenta. In the case of the $\phi$-MHV amplitudes on can avoid all boundary terms and non-factorising 3-point loop amplitudes by choosing $\vec{p}_1 = p_1 + z[2][1]$ and $\vec{p}_2 = p_2 - z[2][1]$. This shift gives us a recursion relation in terms of lower point $\phi$-MHV amplitudes, known finite $\phi$-amplitudes$^5$ and QCD amplitudes leading to a rational term $R^D$. The recursive part of the rational contribution is defined by,
\[
R^D_4 = \sum_{i} A_{L}^{D}(z)R_{R}(z) + R_{L}(z)A_{R}^{D}(z).
\]
(9)
For the $|1\rangle|2\rangle$ shift the contributing diagrams are shown in fig. 2.

![Figure 2: The direct recursive diagrams contributing to $R_4(\phi,1^-,2^-,3^+,4^+)$ with a $|1\rangle|2\rangle$ shift.](image-url)
The final step is to remove overlap terms which appear due to a double counting of poles in $CR_4$. These are computed by evaluating $CR_4(z)/z$ at the poles in recursion:

$$O_4 = \sum_{\alpha} \frac{CR_4(z_\alpha)P_\alpha(z_\alpha)}{P_\alpha(0)} \quad (10)$$

Collecting results for the four gluon case and constructing the Higgs amplitude yields,

$$R_4(H; 1^-, 2^-, 3^+, 4^+) = G(1, 2, 3, 4, \{ \\} \{ \\}) + G(2, 1, 4, 3, \{ \\} \{ \\}) + G(3, 4, 1, 2, \{ \\} \{ \\}) + G(4, 3, 2, 1, \{ \\} \{ \\}) \quad (11)$$

where

$$G(1, 2, 3, 4, \{ \\} \{ \\}) = -\frac{N_F}{96\pi^2} \left[ \frac{(23)^2(34)^3(41)}{(34)[14][34][13][41][3]} - \frac{(23)^2(34)^3}{(34)[14][34][13][41][3]} + \frac{3}{4} \frac{(12)(23)(34)^2}{(34)[41][14][34][13][41][3]} - \frac{(41)[12](34)^2}{s_{341}[41]^2(34)^2} + \frac{2}{s_{314}[34][23]} + \frac{2}{s_{314}[34][12]} \frac{1}{[41][12]} - \frac{1}{4} \left( \frac{34}{[41]} - \frac{(12)}{[34]} \right) \right]^2. \quad (12)$$

5 Conclusions

We have employed four-dimensional unitarity and recursion relations to compute the one-loop corrections to a specific amplitude involving four gluons and a colourless scalar - the Higgs boson. The amplitude presented here may be useful in computing the gluon fusion contamination of the weak boson fusion signal for events containing a Higgs and two jets at the LHC.

Acknowledgements

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References

Early physics with top quarks at the LHC

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CERN,
1211 Genève 23, Switzerland

The ATLAS and CMS experiments are now in their final installation phase and will be soon ready to study the physics of proton-proton collisions at the Large Hadron Collider. The LHC, by producing 2 fb\(^{-1}\) events per second, will provide more than 8 million top events a year at start-up. In this paper, particular emphasis is given to the \( t\bar{t} \) physics studies that can be performed at the beginning of the LHC running, with a limited amount of integrated luminosity (\( \leq 10 \text{ fb}^{-1} \)).

1 Introduction

In the early days of data taking at the LHC, top physics will have a role of primary importance for several reasons. First of all, since top physics allows for precise studies of the Standard Model (SM) and since the determination of the top mass constraints the Higgs mass via radiative corrections. At start-up, already with the first few fb\(^{-1}\) of integrated luminosity and with a non perfectly calibrated detector, a top signal can be clearly separated from the background and the top pair production cross-section can be extracted at better than 20% accuracy and with negligible statistical error. The first measurement of the top mass will provide feedback on the detector performance and top events can be used to understand and calibrate the detector light jet energy scale and the b-tagging. Additionally in scenarios beyond the SM, new particles may decay into top quarks, therefore a detailed study of the top quark properties may provide a hint on new physics. A good understanding of top physics is also essential since top events are a background for many new physics searches.
2 Early selection of top events in the leptons+jets channel

Since top events are so crucial for the initial phase of data taking, it is important to understand how much integrated luminosity is needed to observe the top signal over the background at startup and the effects of a non-perfectly calibrated detector on its observability. A study that uses a very simple selection in the leptons + jets channel, where $t\bar{t} \rightarrow W^+ bW^- \bar{b}$ with a $W$ decaying hadronically and the other leptonically $W \rightarrow e\nu_e (\mu\nu_\mu)$, has been performed by the ATLAS collaboration $^1$. The selection requires 3 jets with transverse momentum $p_T > 40$ GeV/c and one with $p_T > 20$ GeV/c, one isolated lepton with $p_T > 20$ GeV/c and missing transverse energy $E_T > 20$ GeV/c$^2$. In this selection the b-tagging information is deliberately not used since it might not be optimized and calibrated in the initial phase of data taking. The hadronic top is selected as the 3-jet combination with the highest transverse momentum: 2 out of the 3-jets would be resulting from a $W$ decay, therefore only the combinations with a di-jet invariant mass $|m_{jj} - m_W| < 10$ GeV/c$^2$ are kept. Figure 1 shows the expected distribution of the 3-jet invariant mass in a 100 pb$^{-1}$ integrated luminosity sample. The dominant background is the $W$+jets production giving a contribution of the same order as wrongly reconstructed $t\bar{t}$ events. The signal over background ratio is about 0.7 and the relative statistical error is about 10%. Overall the top cross-section could be determined with a total uncertainty of about 20% with few hundred pb$^{-1}$ of integrated luminosity.

3 Top cross-section evaluation

In the leptons + jets channel, a better accuracy on the cross-section can be obtained by refining the selection and in particular by requiring 2 b-tagged jets. To further reduce the background and combinatorics, a converging kinematic fit to $m_W$ can be applied. With 5 fb$^{-1}$ of integrated luminosity, a recent study by the CMS collaboration $^2$ has extracted the $t\bar{t}$ cross-section with the following errors: $\delta \sigma/\sigma = 0.6\%$ (statistical) $\pm 9.2\%$ (systematical) $\pm 0.5\%$ (luminosity).

While the leptons+jet can be considered as the golden channel since the background can be reduced by using simple cuts and the signal will be visible very soon after start-up, promising results have been obtained also in the di-leptonic and fully hadronic channels, where both $W$'s decay either leptonically ($e, \mu$) or hadronically, respectively. A comparison of the performances in the different search channels, as from recent studies by the CMS collaboration $^2$, can be read from Table 3.
Table 1: Breakdown of statistical, systematical and luminosity errors, main background sources, efficiency and signal over background ratio S/B, for the cross-section studies in the lepton+jets, di-leptonic and hadronic channels. The S/B ratio for the lepton+jets channel doesn’t take into account the background from $tt$.

<table>
<thead>
<tr>
<th></th>
<th>syst (%)</th>
<th>stat (%)</th>
<th>lumi (%)</th>
<th>main syst. (%)</th>
<th>main bkg</th>
<th>eff</th>
<th>S/B</th>
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<tr>
<td>$10fb^{-1}$</td>
<td>9.7</td>
<td>0.4</td>
<td>3</td>
<td>$tt$</td>
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<td>6.3</td>
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<td></td>
<td></td>
<td>$b$-tag</td>
<td></td>
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<td></td>
<td>Pile-up</td>
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<td></td>
<td></td>
<td>3.4</td>
<td>3.2</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$10fb^{-1}$</td>
<td>11</td>
<td>0.9</td>
<td>3</td>
<td>$tt$ with $(W \rightarrow \tau\nu_{\tau})$</td>
<td>$(W \rightarrow \tau\nu_{\tau})$</td>
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<td>5</td>
</tr>
<tr>
<td>di-leptonic</td>
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<td></td>
<td></td>
<td>$b$-tag</td>
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<td>4</td>
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<td>$1fb^{-1}$</td>
<td>20</td>
<td>3</td>
<td>5</td>
<td>JES</td>
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<td>11</td>
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<td></td>
<td></td>
<td>$1/9$</td>
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</table>

4 The top mass measurement

In the lepton + jets channel, after an event selection optimised not to bias the mass measurement, different methods have been exploited to extract the top mass ($m_t$). The simplest is to perform a fit to the invariant mass of the 3 jets arising from the hadronic top decay, but this suffers of the impact of poorly reconstructed jets due to effects of FSR and to the semi-leptonic decay of b-quarks. Another method, less affected by systematic errors, reconstructs event by event the entire $tt$ final state via a $\chi^2$ minimisation based on kinematic constraints: the energies of the leptons and jets, the jet directions and the 3 components of the reconstructed neutrino’s are free to vary within their resolutions; $m_t$ is then fitted in slices of $\chi^2$ and is extrapolated from a linear fit to the $m_t$ value corresponding to $\chi^2 = 0$. Alternatively, an event-by-event likelihood method which convolutes the resolution function of the event, or the so called ideogram, with the expected theoretical template can be used. A method which is appealing since it has independent systematic errors, is to select high $p_T$ top pairs with $p_T > 200$ GeV/c; in this case the 2 top quarks tend to be back to back and this can be used to reduce the backgrounds. Since the 3 jets on one hemisphere tend to overlap, the energy in a cone around the candidate top quark has to be collected making the measurement less sensitive to the jet energy calibration. A summary of the different contributions to the error on $m_t$ for the different methods described above can be found in table 4, as from an ATLAS study.

As for the cross-section, $m_t$ can also be extracted from the di-leptonic and the hadronic channels. The di-lepton channel has a clean signature, but 2 neutrino’s need to be reconstructed, this can be done by applying a constrained fit assuming the W mass and two equal masses for the 2 reconstructed top. With an integrated luminosity of 1 fb$^{-1}$, the statistical error on $m_t$ would be of about 1.5 GeV/c$^2$ and the systematical about 4.2 GeV/c$^2$. In the hadronic channel a kinematic fit can be used to reconstruct both top quarks, but the measurement is affected by large QCD backgrounds. With an integrated luminosity of 1 fb$^{-1}$ the statistical error would be of about 0.6 GeV/c$^2$ and the systematical about 4.2 GeV/c$^2$.

5 Searches for new physics

By reconstructing the top mass spectrum in $tt$ lepton+jets events, resonances originated by the decay process $p\bar{p} \rightarrow X \rightarrow tt$ can be observed. From preliminary studies by ATLAS a 1(2) TeV/c$^2$ mass $Z'$ boson produced with a cross-section of 4(3) pb can be observed at about 3$\sigma$ significance with an integrated luminosity of about 5 fb$^{-1}$. 
Table 2: Expected systematical and statistical error contributions to the top mass measurement expressed is $\delta m_t (GeV/c^2)$ for the 3 methods described in the text: the hadronic mass fit, the kinematic fit and the high $p_T$ selection.

<table>
<thead>
<tr>
<th>Component</th>
<th>had. top</th>
<th>kin. fit</th>
<th>high $p_T$</th>
</tr>
</thead>
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<tr>
<td>light jet E scale (1%)</td>
<td>0.2</td>
<td>0.2</td>
<td>-</td>
</tr>
<tr>
<td>b-jet E scale (1%)</td>
<td>0.7</td>
<td>0.7</td>
<td>-</td>
</tr>
<tr>
<td>b-quark fragmentation</td>
<td>0.1</td>
<td>0.1</td>
<td>0.3</td>
</tr>
<tr>
<td>ISR</td>
<td>0.1</td>
<td>0.1</td>
<td>0.1</td>
</tr>
<tr>
<td>FSR</td>
<td>1.0</td>
<td>0.5</td>
<td>0.1</td>
</tr>
<tr>
<td>combinatorial bkg</td>
<td>0.1</td>
<td>0.1</td>
<td>-</td>
</tr>
<tr>
<td>mass rescaling</td>
<td>-</td>
<td>-</td>
<td>0.9</td>
</tr>
<tr>
<td>Underlying event (10%)</td>
<td>-</td>
<td>-</td>
<td>1.3</td>
</tr>
<tr>
<td>total syst.</td>
<td>1.3</td>
<td>0.9</td>
<td>1.6</td>
</tr>
<tr>
<td>stat. err. @10 fb$^{-1}$</td>
<td>0.05</td>
<td>0.1</td>
<td>0.2</td>
</tr>
</tbody>
</table>

Already with 10 fb$^{-1}$ of data, flavour changing neutral currents, which are not allowed at tree level in the SM, can be observed with a sensitivity 2 orders of magnitude better than at Tevatron$^{6,7}$. Finally by studying the double differential angular distribution of $t\bar{t}$ decay products and by comparing the observed values of the spin correlation observables and the SM expectations, the presence of anomalous couplings, Technicolor, spin 0/2 heavy resonances can be observed. With an integrated luminosity of 10 fb$^{-1}$, the spin correlation observables can be extracted with a 3% and 5% statistical and systematical uncertainty, respectively$^{8,9}$.

6 Conclusions

Top physics provides an excellent environment for calibrating the detector and for testing the SM predictions as well as new physics starting from the early days of data taking at the LHC. A large effort has been made by the ATLAS and CMS Collaborations to be ready to analyse the top events from day one, by searching for better selection cuts, improving the generators and systematic errors understanding and exploring alternative analysis methods and decay channels.

References

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5. J. Heininck et al., *Top mass measurement in single leptonic $t\bar{t}$ events*, CMS NOTE 2006-066 (2006)
FORWARD-BACKWARD CHARGE ASYMMETRY IN Z PRODUCTION AT THE LHC.

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We present here a study on the determination of the effective weak mixing angle, $\sin^2 \theta_{\text{eff}}^{\text{lep}}$, from the measurement of the Forward-Backward Asymmetry with a high a statistical precision, $10^{-4}$. To reach such a precision it is necessary to identify the electrons in the forward regions of the ATLAS detector. It is demonstrated that one can reach an electron, jet rejection of more than 100 with an efficiency on electron reconstruction better than 50%, by using a multivariate analysis.

Keywords: LHC, ATLAS, $Z \rightarrow e^+e^-$, Forward-Backward Asymmetry, $\sin^2 \theta_{\text{eff}}^{\text{lep}}$.

1 Introduction

Many measurements of the theoretical electroweak parameters are sensitive to the quantum corrections effect associated to the Higgs boson. The W mass, $M_W$, and the weak mixing angle, $\sin^2 \theta_{\text{eff}}^{\text{lep}}$, are related quadratically to the top mass and logarithmically to the Higgs mass. Therefore the indirect measurement of the Higgs mass from the top mass and the weak mixing angle provides a check of the Standard Model coherence and a validation of any Higgs discovery.

At the LHC, a $Z$ production rate of $\sim 1.5 \times 10^8$ per year at high luminosity of which $\sim 5 \times 10^6$ decay to an electron-positron pair is expected. The determination of the weak mixing angle from the measurement of the forward-backward asymmetry in the $Z \rightarrow e^+e^-$ with a very small statistical error comparable to the LEP one is possible. The electron channel was chosen instead of the muon due to the limited coverage of the muons (up to $|\eta|= 2.7$).
2 Simulation

The signal and background events are generated with the PYTHIA6.2 and the parametrization CTEQ5L of the structure function is used. The events are then fast simulated and reconstructed with the ATLAST, a fast simulation of the ATLAS detector response.

The events $qar{q} \rightarrow Z/\gamma^* \rightarrow e^+e^-$ were generated in various transverse momentum ($p_T$) ranges of the hard scattering matrix element and with a dilepton mass, $m$, greater than 50 GeV.

The main backgrounds to the electron channel are:

- $pp \rightarrow jj$ (QCD): it is the dominant background where each jet simulates an electron. The cross section of this process is greater by several orders of magnitude than the signal one, and dominates at low transverse momentum.

- $pp \rightarrow t\bar{t} \rightarrow e^+e^-$: The top quark decay into the W boson and b quark, followed by the W decay into electron and neutrino ($t \rightarrow Wb, W \rightarrow e\nu$). It has the same signature as the signal one as the two electrons of the final state can simulate the two electrons from Z.

3 Analysis method

The aim of this analysis is the measurement of the forward-backward charge asymmetry in the $Z \rightarrow e^+e^-$ events and its precision. This measurement provides a determination of the weak mixing angle $\sin^2 \theta_{\text{eff}}$ by using the relation\(^1\,^2\):

$$A_{FB} = b(a - \sin^2 \theta_{\text{eff}})$$

(1)

The selection cuts of our analysis in the electron channel requires an electron transverse momentum higher than 20 GeV ($p_T > 20 \text{GeV}$) to simulate the energy threshold of the electron trigger, and a window of 12 GeV around the Z mass, 85.2 GeV < $M_{(e^+e^-)}$ < 97.2 GeV (Z pole). One requires that one of the two electrons lies in the central region ($|\eta| < 2.5$), while the other electron is either in the central region (case 1) or in the forward region (case 2) up to $|\eta| = 4.9$. In the region $2.5 < |\eta| < 3.2$ the calorimeters used are the EMEC and the HEC and for $|\eta| > 3.2$ the forward calorimeter (FCal) is used. Note that we can’t reconstruct the electron track in the forward region ($2.5 < |\eta| < 4.9$) as the tracking system of ATLAS is limited to the region $|\eta| < 2.5$. In addition, the forward calorimeters have a coarser granularity than in the central one and we expect the electron identification in this region to be less performant than in the central one.

At low rapidity $y_{(e^+e^-)}$ most of the events are produced via the annihilation of the sea quark and sea anti-quark, and the probability that the valence quark and the di-electron boost coincide is then lower. The effect on a cut $|y_{(e^+e^-)}| > 1$ is then studied.

We require the missing transverse energy to be less than 20 GeV ($P_T^{\text{miss}} < 20 \text{GeV}$). This cut rejects efficiently the background coming from $pp \rightarrow t\bar{t}$ channel where the top decays semileptonically.

In practice the forward-backward asymmetry is calculated by a counting method of the forward events $N_F$ with $\cos \theta^* > 0$ and backward events with $\cos \theta^* < 0$ ($\theta^*$ is the polar angle of the electron in the Z rest frame):

$$A_{FB} = \frac{N_F - N_B}{N_F + N_B}$$

(2)

As the distribution of the events $N_F$ and $N_B$ follows a binomial distribution, the error on the two quantities can be written as follows: $\sigma_{N_F} = \sigma_{N_B} = \sqrt{N_F N_B} / \sqrt{N_F + N_B}$ and the $A_{FB}$ error is $\sigma_{A_{FB}} = \sqrt{\frac{1 - A_{FB}^2}{N}}$. 
In this analysis, the electron/jet rejection is studied for a fixed electron efficiency (50%). The data are normalized to the integrated luminosity of 100 fb$^{-1}$ (3 years at high luminosity).

4 Results

Table 1 shows the value of the forward-backward asymmetry measurement, its statistical error and the corresponding error on the weak mixing angle $\sin^2 \theta_{\text{eff}}$. When the two electrons are in the central region, we remark that the asymmetry and its error are unchanged with or without the background due to the higher rejection factor in this region. Fig. 1 left shows the variation of asymmetry versus the rapidity of the two electrons. It is observed that the asymmetry increases by a factor 2 when allowing the second electron to be in $|\eta| < 4.9$. As shown in the right plot of Fig. 1, the accuracy on the forward backward asymmetry improves while the jet rejection increases in the forward regions and it is almost constant for rejection greater than 100. The statistical error reached here (for a forward rejection of 100 and forward electron efficiency of 50%) on the weak mixing angle is of $\sim 10^{-4}$.

The statistical error on $\sin^2 \theta_{\text{eff}}$ is deduced from the relation 3 where the parameters, a and b, are derived from theory$^2$ including the radiative corrections.

<table>
<thead>
<tr>
<th>Rej$gy$</th>
<th>$A_{FB}$ (%)</th>
<th>$\delta A_{FB}$</th>
<th>$\delta \sin^2 \theta_{\text{eff}}$</th>
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<tr>
<td>$</td>
<td>y_g^-</td>
<td>,</td>
<td>y_g^+</td>
</tr>
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<td>y_g^+</td>
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<td>y_Z</td>
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</tr>
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<td>y_g^+</td>
<td>&lt; 4.9$</td>
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<tr>
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<td>$1.08 \times 10^{-4}$</td>
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<td>10 1.52</td>
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<td>1 0.43</td>
<td>$3.47 \times 10^{-4}$</td>
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Table 1: Values of the FB asymmetry, of the statistical error on the asymmetry and on the weak mixing angle for rejection values in the forward region ($10^4$, 100, 10 and 1).

5 Forward electron reconstruction

In order to evaluate our analysis result and demonstrate if we can reach such requirements defined by our analysis, a multivariate analysis was used to evaluate the performance of the forward calorimeters in the separation of electrons from hadrons. The input variables used describe the shower shape (lateral and longitudinal) and its development. The most discriminating variable is the fraction of the total energy deposited in the most energetic cell. Various discriminant methods are used to confirm the obtained result. Fig. 2 shows the variation of the jet rejection versus electron efficiency in the forward regions for three different analysis methods (Fisher
Figure 1: Left: Forward-Backward asymmetry versus dilepton rapidity in the case 1 (red points) and in the case 2 (black points). Right: Forward-Backward asymmetry accuracy versus the forward electron/jet rejection in the events of the case 2.

We found that a rejection of 100 can easily be obtained while keeping an electron efficiency better than the 50%.

Figure 2: Background rejection versus the signal efficiency in the EMEC (left) and in the forward calorimeter (right).

References

IX - Structure Functions, Spin and Diffraction Session

Chairperson: Urs Wiedemann
MEASUREMENT OF THE GLUON POLARISATION AT COMPASS

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COMPASS measurements of the gluon polarisation in nucleon, $\Delta G/G$, are reviewed. Two different approaches based on tagging the Photon Gluon Fusion process are described. They rely on the open charm meson or high-$p_T$ hadron pairs detection. The obtained results are:

$-0.57 \pm 0.41$ (stat.) $\pm 0.17$ (syst.) for the open charm, $0.06 \pm 0.31$ (stat.) $\pm 0.06$ (syst.) and $0.016 \pm 0.058$ (stat.) $\pm 0.055$ (syst.) for high-$p_T$ for $Q^2 < 1$ (GeV/$c$)$^2$ and $Q^2 > 1$ (GeV/$c$)$^2$

regimes, respectively.

1 Introduction

In the framework of QCD the nucleon spin can be decomposed into four contributions: from quarks $\Delta \Sigma$, from gluons $\Delta G$ and from angular momenta of quarks and gluons $L_q$ and $L_g$:

$$\frac{1}{2} = \frac{1}{2}\Delta \Sigma(\mu^2) + \Delta G(\mu^2) + L_q(\mu^2) + L_g(\mu^2)$$

(1)

where $\mu^2$ is a scale at which the nucleon is probed. $\Delta \Sigma$ is determined precisely in a QCD fit to the $g_1$ structure function data and is $0.30 \pm 0.01$ (stat.) $\pm 0.02$ (syst.) at $3$ (GeV/$c$)$^2$.

This method also gives $\Delta G/G$ albeit with large uncertainty due to the limited kinematical range of the $g_1$ measurements. One of the goals of COMPASS is a determination of the $\Delta G/G$ quantity. The method relies on tagging the Photon Gluon Fusion (PGF) process with high-$p_T$ hadron pairs, a channel studied also by HERMES$^{2,3}$ and SMC$^4$. In COMPASS also a direct channel based on the PGF tagging with the open charm meson production and decay is used.

The COrmon Muon and Proton Apparatus for Structure and Spectroscopy (COMPASS) is a two-stage magnetic spectrometer located at CERN SPS muon beam line which delivers 160 GeV/$c$ positive muons of intensity $2 \cdot 10^8$ particles per 16.8s SPS cycle. Muons in the beam are naturally polarised. The polarisation is -0.76 (for 2002 and 2003) and -0.81 (for 2004 data taking period). The COMPASS target is composed of two solid state $^6$LiD cells, each 60 cm long and 3
cm in diameter. The spins of the deuterons in the cells are polarised in the opposite directions, parallel and anti-parallel to the beam polarisation. The \( ^6\text{Li} \) basically consists of a deuteron plus \(^4\text{He} \) core; a dilution factor of about 0.4 is obtained for the target material. The average polarisation of the target deuterons is 50%. Directions of polarisation are flipped every 8 hours. The twin-target and the spin reversals are needed to cancel false asymmetries originating from different spectrometer acceptances for the two target halves and from the time variations of the beam flux. Particles produced in the interactions are traced and identified in two spectrometers equipped with tracking and identification detectors (including RICH), magnets and calorimetry\(^5\).

2 Determination of \( \Delta G/G \)

The analysis is based on the measurement of the cross sections asymmetry of the Photon Gluon Fusion (PGF) interactions with different relative spin orientations of the projectile and of the target nucleon. This asymmetry, \( A^{\gamma N} \), is coupled to \( \Delta G/G \) via:

\[
A^{\gamma N} = R_{\text{PGF}} \hat{a}_{\text{LL}} \frac{\Delta G}{G} + A_{\text{BG}}
\]

where \( \hat{a}_{\text{LL}} \) is the partonic asymmetry, \( R_{\text{PGF}} \) is the fraction of PGF events in the selected sample and \( A_{\text{BG}} \) is the background asymmetry. Two methods to extract \( \Delta G/G \) are discussed below.

2.1 The open charm method

The mass of a charm quark is much larger than that of \( u, d \) and \( s \) quarks. The intrinsic charm content of the nucleon at COMPASS kinematics is negligible. Also the production of the charm quarks in the fragmentation process is highly suppressed. The only reaction that has a significant contribution to the charm production is PGF process. The charmed quarks fragment subsequently into charmed hadrons. The detection and identification of them provides a clear tag for the PGF. The studies rely on the detection of \( D^0 \) mesons in their \( \pi K \) decay channel. The mesons are reconstructed by pairing each two charged tracks from a given event and calculating the invariant mass of the system. The signal-to-background ratio is small, of the order of 1:10. Therefore, a second, more exclusive decay channel is studied: \( D^* \rightarrow D^0 \pi \rightarrow K \pi \pi \). In this channel a cut on the \( D^* \) mass is imposed and the signal-to-background ratio is increased to approximately 1:1.

Each open charm event is characterised by different value of \( \hat{a}_{\text{LL}} \). The \( R_{\text{PGF}} \), corresponding in the open charm channel to the signal-to-background ratio, is a function of the \( K \pi \) invariant mass. Therefore a weighted method of the \( \Delta G/G \) extraction is applied. Each event is weighted with its \( \hat{a}_{\text{LL}} \) and \( R_{\text{PGF}} \). This requires knowledge of the \( \hat{a}_{\text{LL}} \) on an event-by-event bases. As in the analysis only one charmed meson is required, the reaction kinematics is unknown and the \( \hat{a}_{\text{LL}} \) cannot be calculated. Thus the parameterisation based on the measured kinematical variables is introduced providing the estimation of the \( \hat{a}_{\text{LL}} \) for each event. The parameterisation is obtained using Neural Networks trained on the sample prepared with AROMA generator in LO QCD. The correlation between the generated and reconstructed \( \hat{a}_{\text{LL}} \) is 82%.

The major contributions to the systematic error are listed in Tab 1. The background asymmetry is checked in the signal sidebands, in the wrong charge combinations and by simultaneous fitting the signal and the background asymmetries. The data is divided into subsamples, each containing events recorded in approximate a week long intervals. For each the \( \Delta G/G \) is calculated. Dispersion of the values is used to estimate the false asymmetries arising from the detector instabilities. The parameter sets in AROMA are varied and a number of \( \hat{a}_{\text{LL}} \) parametrisations is prepared and used to check the stability of the obtained \( \Delta G/G \). Around 300 different fits are used for fitting the signal and the background and the calculation of the \( R_{\text{PGF}} \).
Table 1: Major contributions to the systematic uncertainty. They are estimated independently for both channels and found to be equal. The contribution from the uncertainties of the beam and target polarizations, dilution factor and the correlation between the signal strength and the \( \delta_{LL} \) added in quadratures are presented as "Other".

<table>
<thead>
<tr>
<th>Background asymmetry</th>
<th>0.07</th>
<th>Fitting</th>
<th>0.09</th>
</tr>
</thead>
<tbody>
<tr>
<td>False asymmetries</td>
<td>0.10</td>
<td>Other</td>
<td>0.07</td>
</tr>
<tr>
<td>Parameters of AROMA</td>
<td>0.05</td>
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</tr>
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A preliminary result from the open charm analysis for 2002-2004 data is:

\[
\langle \frac{\Delta G}{G} \rangle = -0.57 \pm 0.41 \text{ (stat.)} \pm 0.17 \text{ (syst.)}
\]  \hspace{1cm} (3)

The average fraction of the nucleon momentum carried by the gluon, \( x_g \), for the selected sample is 0.15 with RMS of 0.08. The \( \mu^2 \) is given by the mass of the charm quark and the \( D^0 \) meson transverse momentum with respect to photon: \( \mu^2 = 4(m_c^2 + p_T^2) \) and equal 13 (GeV/c)^2.

2.2 The high-\( p_T \) method

The high-\( p_T \) method relies on a sample of PGF events with light quark pair production. With the selection of two high-\( p_T \) hadrons a fraction of PGF events is enhanced and reaches 30%. The high-\( p_T \) analysis is performed in two kinematical regimes: \( Q^2 > 1 \) (GeV/c)^2 and \( Q^2 < 1 \) (GeV/c)^2.

For the high-\( p_T \) analysis with \( Q^2 > 1 \) (GeV/c)^2 the sample contains a large fraction of the Leading Order (\( \gamma q \rightarrow q \)) and QCD-Compton interactions (\( \gamma q \rightarrow qg \)) apart from the PGF. The fraction of the resolved photon processes is assumed to be negligible. The cut \( x < 0.05 \) suppresses the Leading Order and QCD-Compton contributions. Thus:

\[
A_{\gamma^* N \rightarrow h h} = R_{PGF} \langle \Delta G \rangle \frac{\Delta G}{G}
\]  \hspace{1cm} (4)

The ratio \( R_{PGF} \) is estimated from MC simulation using the LEPTO 6.5.1 generator.

About 90% of events containing two high-\( p_T \) hadrons have \( Q^2 < 1 \) (GeV/c)^2. Here apart from the PGF, Leading Order and QCD-Compton the resolved photon processes have to be taken into account. They contribute to more than 50% of interactions. Three channels involving resolved photon dominate: \( qg^* \rightarrow qq, \; qg^* \rightarrow qg, \; gg^* \rightarrow gg \). The total asymmetry is:
\[ A^{\gamma N \rightarrow hh} = \left[ R_{PQGF}(\bar{a}_{LL}^{QGF}) + R_{gq}(\bar{a}_{LL}^{gq}) \left( \frac{\Delta G}{G} \right)^\gamma \right] \frac{\Delta G}{G} + \]
\[ + \left[ R_{LO}(\bar{a}_{LL}^{LO}) + R_{QCD}(\bar{a}_{LL}^{QCD}) + R_{qg}(\bar{a}_{LL}^{qg}) \left( \frac{\Delta G}{G} \right)^\gamma + R_{gq}(\bar{a}_{LL}^{gq}) \left( \frac{\Delta G}{G} \right)^\gamma \right] \frac{\Delta q}{G} \]

where the superscript \( \gamma \) denotes the parton distributions describing the resolved photon structure. The fractions of the processes as well as the average \( \bar{a}_{LL} \) for each of them are estimated with a MC simulation using PYTHIA 6.2. The large number of parameters not measured directly but estimated from the simulation results in a model dependence of the resulting \( \Delta G/G \) value. This dependence is encompassed in a systematic error (see also \( ^6 \)).

The preliminary results from the high-\( p_T \) analysis from the 2002-2003 data at \( Q^2 > 1 \) (GeV/c)^2 and for 2002-2004 data at \( Q^2 < 1 \) (GeV/c)^2 are:

\[ \langle \frac{\Delta G}{G} \rangle = 0.06 \pm 0.31 \text{ (stat.)} \pm 0.06 \text{ (syst.)} \] (6)
\[ \langle \frac{\Delta G}{G} \rangle = 0.016 \pm 0.058 \text{ (stat.)} \pm 0.055 \text{ (syst.)} \] (7)

The average \( x_g \) for the first one is 0.13 with an RMS of 0.08 while for the second one it is 0.095 \( ^{+0.08}_{-0.04} \). The scale \( \mu^2 \), given by the transverse momentum of the outgoing partons with respect to photon, is 3 (GeV/c)^2 for both sets.

In Fig. 2 these direct measurements are compared to parameterisation of \( \Delta G(x_g) \) obtained in a NLO fit to the \( g_1 \) data including the new deuteron results \( ^1 \). Two solutions for \( \Delta G(x_g) \) were found, one positive and one negative, both resulting in small \(|\Delta G| \sim 0.2 - 0.3\) for \( \mu^2 = 3 \) (GeV/c)^2. The error band corresponds to the change of \( \chi^2 \) by unity. With the present precision measurements cannot distinguish between two possible scenarios. However they are in line with a small value of \( \Delta G \).

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LONGITUDINAL DOUBLE SPIN ASYMMETRY IN JET PRODUCTION AT STAR

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We report measurements of the longitudinal double-spin asymmetry $A_{LL}$ for the inclusive production of jets at midrapidity in polarized proton-proton collisions at $\sqrt{s} = 200\text{GeV}$. The data amount to an integrated luminosity of $3\text{pb}^{-1}$ and were collected with the STAR detector at the Relativistic Heavy Ion Collider. Typical beam polarizations were 50%. The $A_{LL}$ measurements cover jet transverse momenta $5 < p_T < 30\text{ GeV}/c$ and provide sensitive constraints on the gluon spin contribution to the proton spin.

1 Introduction

The Relativistic Heavy Ion Collider (RHIC) is the first polarized high-energy proton-proton collider, providing collisions at $\sqrt{s} = 200\text{GeV}$ and in the future at $\sqrt{s} = 500\text{GeV}$. One of the main objectives of the program with polarized proton beams is the measurement of the gluon spin contribution, $\Delta G$, to the proton spin$^1$. The processes under study include inclusive jet and pion production, dijets, heavy flavor, and photon-jet coincidences. The present focus is on inclusive measurements with large production cross sections. At a center-of-mass energy $\sqrt{s} = 200\text{GeV}$ production processes at midrapidity and with transverse momenta $p_T > 5\text{GeV}/c$ typically give sensitivity to the integral of $\Delta G$ over the range $0.03 \lesssim x \lesssim 0.3$ in the gluon momentum fraction. Coincidence measurements can provide the $x$ dependence of $\Delta G$. However, they require larger total integrated luminosities than have been sampled so far. Collisions at $\sqrt{s} = 500\text{GeV}$ will extend the $x$ coverage to smaller values, which are important for global analyses of the polarized parton distributions. The STAR (Solenoidal Tracker at RHIC) detector$^2$, with its large acceptance and electromagnetic calorimetry, is well suited for jet reconstruction at RHIC.

The (polarized) jet cross section has large contributions from (polarized) gluon-gluon and quark-gluon scattering, which provide direct sensitivity to $\Delta G$ in the measurements of the longitudinal double-spin asymmetry $A_{LL}$,

$$A_{LL} = \frac{\sigma^{++} - \sigma^{+-}}{\sigma^{++} + \sigma^{+-}},$$  \hspace{1cm} (1)

where $\sigma^{++}$ ($\sigma^{+-}$) is the inclusive jet cross section when the colliding proton beams have equal (opposite) helicities. The gluon-gluon scattering contribution dominates up to jet $p_T \lesssim 8\text{GeV}/c$ while at higher $p_T$ in the range of the present measurements the quark-gluon scattering contribution is the largest. The quark-quark contribution is small$^3$.

First inclusive jet $A_{LL}$ and cross section results from STAR have been published$^4$. The unpolarized cross section is well described by NLO pQCD calculations for jet $p_T$, $5 < p_T < 50\text{GeV}/c$,.
motivating the application of the NLO pQCD framework to interpret the spin asymmetry results. The asymmetry $A_{LL}$ was measured for $5 < p_T < 17$ GeV/c and disfavors large positive gluon polarization in the proton. These proceedings report on new measurements of $A_{LL}$ for inclusive jet production with increased sensitivity and extended coverage in jet $p_T$.

2 Experiment and Data Analysis

The present results are based on an integrated luminosity of $\sim 3\text{pb}^{-1}$ and were recorded in the year 2005. The STAR detector subsystems used for jet reconstruction are the Time Projection Chamber (TPC) and the Barrel Electromagnetic Calorimeter (BEMC). The TPC provides the momentum of charged particles in the pseudorapidity range $-1.3 \leq \eta \leq 1.3$ for all azimuthal angles $\phi$. The BEMC is a lead-scintillator sampling calorimeter which measures electromagnetic energy deposits. During the data taking period in 2005, the BEMC covered $0 < \eta < 1$ and $0 < \phi < 2\pi$. The remaining half, covering $-1 < \eta < 0$, was commissioned before the data taking in 2006. Beam-Beam Counters (BBC), located on each side of the STAR interaction region, provided the beam collision trigger and were used to measure the relative luminosities for different helicity configurations $^5$ as well as transverse beam polarization components at STAR.

The majority of the jet data were collected using a new dedicated jet-patch (JP) trigger that required a transverse energy sum in at least one of six BEMC patches, each covering $\Delta \eta \times \Delta \phi = 1 \times 1$. The JP trigger efficiency is higher than the efficiency of the high tower trigger (HT), which selects on an energy deposit in a BEMC tower of size $\Delta \eta \times \Delta \phi = 0.05 \times 0.05$ and was used also in previous data taking periods. In addition, the JP trigger selects a less biased distribution of jets than the HT trigger does. The HT and JP triggers were used with two different energy thresholds.

Jets were reconstructed with a midpoint cone algorithm $^6$ with a cone radius of 0.4 and using charged TPC tracks and BEMC energy deposits. The details for other parameters can be found in Ref. $^4$. Only jets with reconstructed transverse momenta $p_T > 5$ GeV/c that fulfilled the trigger conditions were considered in the analysis. Jets with their axis between a nominal $\eta$ of 0.2 and 0.8 were selected so that the effects from the BEMC acceptance edges were small. BBC timing information was used to accept events with vertex positions along the beam direction in the inner region of STAR for which tracking efficiencies are uniform. The same timing information was used in the beam luminosity measurement. Beam background caused an occasional signal in the BEMC. Its contribution to the jet yield was suppressed by requiring the ratio of jet energy in the BEMC to the total jet energy to be between 0.1 and 0.8.

3 Results

The jet yields were sorted by equal $(N^{++})$ and opposite $(N^{+-})$ helicity combinations of the colliding proton beams and $A_{LL}$ was evaluated according to:

$$A_{LL} = \frac{\sum (P_1 P_2)(N^{++} - R N^{+-})}{\sum (P_1 P_2)^2(N^{++} + R N^{+-})},$$

where $P_1 P_2$ is the product of the measured beam polarizations $^7,8$, and $R$ is the measured ratio of luminosities for equal and opposite proton beam helicities $^9$. The average online beam polarization was $\sim 50\%$. The ratio $R$ was between 0.8 and 1.2, and was measured to $O(10^{-3})$ accuracy. The yields $N^{++}$ and $N^{+-}$ were recorded concurrently since the proton beam helicities alternated for successive beam bunches in one beam and for successive pairs of beam bunches in the other beam. To further minimize systematic effects in the measurement of $A_{LL}$, the beam helicity pattern was alternated between RHIC beam fills.
The asymmetry was calculated for the combined sample of about 1.97M jets of JP and HT triggered data. Figure 1 shows $A_{LL}$ as a function of jet $p_T$ with statistical error bars for the present and published data, and the systematic uncertainty band for the present data, not including the 25% scale uncertainty from the online beam polarization measurement. The leading contributions to the systematic uncertainty band arise from the bias introduced by the BEMC trigger requirements and from a conservative upper limit on possible false asymmetries in the measurement. Other systematic uncertainties include effects from non-longitudinal beam polarization components at the STAR interaction region, from uncertainty in the measurement of $R$, and from beam background. Systematic checks with randomized beam-spin patterns showed no evidence for bunch-to-bunch or fill-to-fill systematics in $A_{LL}$.

The curves in Figure 1 show NLO pQCD evaluations for $A_{LL}$ based on commonly used polarized parton distribution functions $^{3,10}$. The curve labeled GRSV-std is based on a best fit to inclusive DIS data $^{10}$. The other curves correspond to maximally positive ($\Delta g = g$), negative ($\Delta g = -g$), or vanishing ($\Delta g = 0$), gluon polarizations at a 0.4 GeV$^2$/c$^2$ initial scale of the parton parametrizations $^{10}$.

The present data are in good agreement with the our published results $^4$ in the region of kinematic overlap $5 < p_T < 17$ GeV/c, where they improve significantly the precision, and extend jet $p_T$ up to 30 GeV/c. The corresponding range of gluon momentum fractions sampled by data is $0.03 < x < 0.3$. The fraction of the first moment of the GRSV-std polarized gluon distribution is about half over this $x$ range. The data are not consistent with the GRSV scenario of maximal positive gluon spin contribution to the proton spin ($\Delta g = g$).

An additional 8.5 pb$^{-1}$ was sampled in 2006 with beam polarizations of $\sim 60\%$. This will improve the precision of inclusive $A_{LL}$ measurements and offers good prospects for dijet analyses.
4 Summary

We reported preliminary measurements of the longitudinal double-spin asymmetry $A_{LL}$ for the inclusive production of jets in polarized proton-proton collisions at $\sqrt{s} = 200$ GeV with jet $p_T$ up to 30 GeV/c. The $A_{LL}$ data, compared to the commonly used GRSV set of polarized parton distributions, exclude the scenario with a large positive gluon spin contribution to the nucleon spin. Future inclusion of our data as well as data in Ref. 11,12 in global analyses should improve our knowledge of the polarized gluon distribution for $0.03 \lesssim x \lesssim 0.3$. First promising works in this direction have already been published13.

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Measurement of Transverse Spin Effects at COMPASS

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By measuring transverse single spin asymmetries one has access to the transversity distribution function $\Delta_T q(x)$ and the transverse momentum dependent Sivers function $q_T^J(x, k_T)$. New measurements from identified hadrons and hadron pairs, produced in deep inelastic scattering of a transversely polarized $^4$LiD target are presented. The data were taken in 2003 and 2004 by the COMPASS collaboration using the muon beam of the CERN SPS at 160 GeV/c, resulting in small asymmetries.

Keywords: transversity, transverse spin asymmetry, Collins asymmetry, Sivers asymmetry, COMPASS

1 Introduction

Measurements\(^1,2\) showed, that the effects of transverse spin in high energy hadronic physics are not naturally suppressed, as it was assumed\(^3\). On the contrary, transverse spin asymmetries provide a way for a measurement of transversity\(^4\) and the quark transverse momentum $k_T$ dependent distribution function $q_T^J(x, k_T)$, the distribution of unpolarized quarks in a transversely polarized nucleon\(^5\). Denoting the Bjoerken scaling variable by $x$, the spin structure of the nucleon can be described at leading order by three leading twist distribution functions, the unpolarized quark distribution $q(x)$, the helicity distribution function $\Delta q(x)$, and the transversity $\Delta_T q(x)$. In the quark parton model (QPM) $\Delta_T q(x)$ can be interpreted as the difference of distributions of quarks polarized parallel and antiparallel to the nucleon spin in a transversely polarized nucleon. Thus $\Delta q(x)$ and $\Delta_T q(x)$ are identical in a non-relativistic picture of the nucleon. But in the QPM $\Delta q(x)$ can be thought of as describing the nucleon spin structure in a frame that is boosted parallel to the nucleon spin, whereas $\Delta_T q(x)$ in a frame that is boosted in the transverse direction. Since rotational symmetry is broken under boosts, $\Delta_T q(x)$ provides complementary information to the proton spin puzzle. The three distribution functions are of equal importance, however, in comparison to the first two, knowledge of the transversity is quite scarce, since it remains inaccessible in inclusive DIS measurements due to its chiral odd nature.

However, semi inclusive measurements, where at least one hadron fragmenting via a chiral odd fragmentation function in the final state is detected, allow to probe transversity. Because the product of a distribution function and fragmentation function is again chiral even, it can be observed in transverse single spin asymmetries. The chiral odd fragmentation of a transversely polarized quark into a unpolarized hadron can be described by the Collins fragmentation function $\Delta_T^D D_q^h$ and the fragmentation into two unpolarized hadrons by the two hadron interference fragmentation function $H_\perp^T$. Denoting the unpolarized fragmentation function by $D_q^h$, the spin dependent fragmentation of a quark can be written as $D_q^h + \Delta_T^D D_q^h \sin \Phi_C$, if one hadron is pro-
Figure 1: Coordinate systems for azimuthal angles. The x-z plane is defined by the incoming and scattered muon. The z-axis is the virtual photon. In the one hadron case (left) the initial state and final state polarization vectors are denoted s and s'.

Produced and $D_q^h + H_1^s \sin \Phi_R^{11, 12, 13, 14}$ if two hadrons are produced. $\Phi_C$, the Collins angle and $\Phi_R$ are the azimuthal angles between the polarization vector of the fragmenting quark and the momentum of the produced hadron and the vector $\vec{R}$, describing the two hadron system, respectively. The relevant coordinate systems are depicted in fig. 1. $\vec{R}$ is the linear combination of the momenta of the produced hadrons, weighted by the relative energy transfer from the scattered muon, to achieve a definition of azimuthal angles that is invariant against boosts along the photon direction: $\vec{R} = \frac{(2z_1 - x_1)^2}{z_1 + z_2}$. By measuring the angular dependence of the produced hadrons on $\Phi_C$ and $\Phi_R$, respectively, it is possible to probe the transversity distribution. For the Collins asymmetry one relies on the measurement of transverse momentum originating from the fragmentation process. However, the azimuthal dependence on $\Phi_R$ of the cross section for the process $\Delta_Tq \otimes H_1^s$ should remain after integrating out transverse momenta. If one leaves the collinear picture and allows transverse momenta of the quarks, more distribution functions of quarks exist. One of these, the so called Sivers function $q_0^T(x, \vec{k_T})$, describes the distribution of unpolarized quarks in a transversely polarized nucleon. It is strongly connected to the angular momenta of quarks in the nucleon, which might be another contribution to the nucleon spin. The Sivers function can be probed via the Sivers effect, where the correlation of the quark transverse momentum with the nucleon spin leads to the dependence of the SIDIS cross section on the azimuthal angle between the nucleon polarization vector and the momentum of the produced hadron. This angle $\Phi_S$ is called Sivers angle. Since the angular dependence of the cross section on $\Phi_S$ and $\Phi_C$ are orthogonal functions, Sivers and Collins effects can be disentangled with a transversely polarized target.

2 Experimental results

COMPASS is a fixed target experiment with a broad physics program at the M2 beamline at CERN and is described in detail in ref. 6. For about 20% of the data taken in the years 2002, 2003 and 2004 a transversely polarized $^6$LiD target is used, which has a favourable dilution factor of $f \approx 0.4$ and a polarization of about 50%. The target consisted of two cells which were polarized in opposite directions and polarization was reversed every 4-5 days to minimize systematic effects. A Ring Imaging Cherenkov (RICH-1) detector and two hadron calorimeters provide particle identification capabilities. RICH-1 is a gas RICH with a 3m long $CaF_2$ radiator. It is characterized by large transverse dimensions in order to cover the whole spectrometer acceptance ($\pm250 \times 200$ mrad) and was operational during data taking in the years 2003 and 2004. Asymmetries for unidentified hadrons were already published. For the analysis presented, events with an incoming beam track crossing both target cells, a scattered muon track and at least one outgoing hadron for Collins and Sivers asymmetry extraction or two outgoing
hadrons with opposite charge for the two hadron correlation extraction, are selected. Positive identification of the hadrons in the final data sample by RICH-1 was required. Clean hadron and muon selection was achieved using the hadron calorimeters and considering the amount of traversed material. To select DIS events, cuts on $Q^2 > 1({\text{GeV/c}})^2$ and $W > 5\text{GeV}/c^2$ were made, $Q^2$ being the photon virtuality and $W$ the mass of the final hadronic state. Additional cuts are applied to ensure that from the hadron sample the relevant physics signal can be extracted. Requiring the relative energy in the muon scattering process $y$ to fulfill $0.1 < y < 0.9$ limits the error due to radiative corrections (higher cut) but warrants that the energy loss of the scattered beam particle is high enough to allow for reliable tracking (lower cut). For the one hadron asymmetries a lower limit of 0.2 for the relative energy $z$ of the hadron is demanded. The underlying reasoning is that in the string fragmentation process hadrons with a higher energy are more sensitive to the properties of the struck quark spin. For the two hadron correlation the cut is $z_1, z_2 > 0.1$ and in addition $z_1 + z_2 < 0.9$ to avoid the kinematic region of exclusive $\rho$ production. After all the cuts, $5.3 \cdot 10^6$ positive pions, $4.6 \cdot 10^6$ negative pions and $9.5 \cdot 10^5$ positive kaons and $6.2 \cdot 10^5$ negative kaons remain for the single hadron analysis. The two hadron correlation signal is extracted from $3.7 \cdot 10^6 \pi^+ - \pi^-$, $2.4 \cdot 10^5 \pi^+ - K^-$, $3.0 \cdot 10^8 K^+ - \pi^-$ and $8.6 \cdot 10^4 K^+ - K^-$ pairs. From these samples the respective asymmetries $A_j$ are extracted by azimuthal count rate asymmetries for target cells with different polarizations. The count rate in the upstream and downstream target cell (k=u,d) for the two polarisations (+,−) $N^\pm_{j,k}$ in a given $\Phi_j$ bin (j=C,S,R) can be written as

$$N^\pm_{j,k} = F_k^{\pm} n_k \sigma_{\Phi_j}^{\pm}(\Phi_j) \cdot (1 \pm \epsilon_{j,k}^\pm \sin \Phi_j)$$

(1)

Here $F$ is the muon flux, $n$ the number of target particles, $\sigma$ the spin averaged cross section, $a_j$ the product of the angular acceptance and efficiency of the spectrometer. The quantity $\epsilon_{j,k}^\pm$ is $f_j \cdot |P_{T,j,k}^\pm| \cdot D_{NN} \cdot A_j$. Where $f$ is the dilution and $|P_{T,j,k}^\pm|$ the polarisation of the target, $D_{NN}$ the spin transfer of the virtual photon to the fragmenting quark and $A_j$ the asymmetry. For Collins and two hadron asymmetries $D_{NN}$ can be calculated in QED as $D_{NN} = \frac{1-y+\lambda/2}{y}$. Since the Sivers effect probes unpolarized quarks $D_{NN}$ does not enter into the corresponding asymmetry. With the obtained count rates the double ratio product

$$A_j(\Phi_j) = \frac{N_{j,u}^+(\Phi_j)}{N_{j,u}^-(\Phi_j)} \cdot \frac{N_{j,d}^+(\Phi_j)}{N_{j,d}^-(\Phi_j)}$$

(2)

is build and fitted with $p_0 \cdot (1 + A_j^m \sin \Phi_j)$. The raw asymmetry $A_j^m = \epsilon_{j,u}^+ + \epsilon_{j,d}^+ + \epsilon_{j,u}^- + \epsilon_{j,d}^- = 4 \epsilon_{j,u}$ is calculated for kinematical binning in $z$, the relative energy of the produced hadron, $p_t$, its transverse momentum and $x$, the Bjorken scaling variable in the one hadron case, whereas it is built for bin in $z$, $M_{inv}$ and $x$ in the two hadron case. $z$ and $M_{inv}$ denote the relative energy and invariant mass of the two hadron system. The corrected asymmetries are shown in figures 2 and 3. Phenomenological work on the Sivers effect has shown that HERMES results for protons and COMPASS results on deuteron may be described within the same theoretical frame, at least at the present level of accuracy of the data. The obtained asymmetries are small and agree well with model calculations, that predict suppressed signals due to the isoscalar target. A model of transversity from the chiral quark soliton model and Collins fragmentation function extracted from a fit to HERMES proton data shows that the favoured and unfavoured Collins fragmentation function seem to be of the same magnitude but of opposite sign. The predictions obtained from this model agree well with the measured asymmetries for unidentified hadrons. Similarly the predictions for the two hadron asymmetries depending on the convolution of transversity with $H_1^c$ are predicted to be small. Due to an interference term in the two pion production one model predicts a strong dependence on the invariant mass around the $\rho$ mass, which cannot be observed in the current COMPASS data. Measurements taken so far on a
deuteron target allow constraints on models for the d-quark Sivers and transversity distribution. They also point to the absence of a gluon contribution to the orbital angular momentum of the partons in the nucleon.

3 Outlook

COMPASS continues data taking in 2007 with a transversely polarized proton target. In combination with the data already measured on deuteron flavour separation for transversity and Sivers distribution function will be possible. An additional analysis is planned for leading hadron pair asymmetries

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PROTON STRUCTURE FUNCTIONS AT HIGH $Q^2$ AND HIGH $x$ AT HERA

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Proton structure measurements at high $Q^2$ and high $x$, performed by the H1 and ZEUS collaborations at the HERA collider, are reviewed. Neutral and charged current deep inelastic scattering cross sections and structure functions are presented. The review also discusses improvements to the parton density measurements using jet cross section data and recent high $Q^2$ inclusive cross section measurements. The projected parton density uncertainties using the entire HERA data set are also presented.

1 Introduction

Precise measurements of the proton parton density functions (PDFs) are crucial for understanding the structure of the proton. This is of particular importance with the imminent start of the LHC proton-proton collider. Unpolarised lepton beams were used before the luminosity upgrade in 2000 (HERA-I), whereas the post-upgrade collider (HERA-II) has delivered polarised lepton beams. This paper reviews the latest measurements at high $Q^2$ and high $x$ performed by the ZEUS and H1 collaborations at HERA.

The PDFs are determined in global fits at next-to-leading order (NLO) in QCD using data from deep inelastic scattering (DIS) experiments. The kinematic range covered by HERA has allowed the determination of PDFs across a wide range of phase space spanned by the fractional proton momentum of the struck quark, Bjorken-$x$, and the negative squared four-momentum transfer, $Q^2$, of approximately $10^{-3} < x < 1$ and $0 < Q^2 < 10^5$ GeV$^2$.

2 Cross Section Measurements and Structure Functions

The Born-level reduced cross section for the $e^+p$ neutral current (NC) interaction with polarised lepton beams can be written as$^1$,

$$\bar{\sigma}^{e+p} = \frac{xQ^4}{2\pi\alpha^2 Y_+} d^2\sigma[\bar{e}^+p] = F_2^\pm(x, Q^2) + \frac{Y_-}{Y_+}xF_3^\pm(x, Q^2) - \frac{y^2}{Y_+}F_L^\pm(x, Q^2),$$

where $\alpha$ is the fine-structure constant, $Y_+ \equiv 1 \pm (1 - y)^2$ and $y$ is related to the centre-of-mass energy, $\sqrt{s}$, via $Q^2 = sxy$. The longitudinal structure function, $F_L$, is small in the kinematic region considered and can be ignored. The structure functions, $F_2$ and $xF_3$, contain the sum and difference of the quark and anti-quark PDFs and can be separated into contributions from pure $\gamma$ exchange, the interference of $\gamma$ and $Z$ boson exchange and from pure $Z$ exchange. These
terms depend on the lepton beam charge, the longitudinal polarisation of the lepton beam, \( P_e \), the mass of the Z and W bosons, \( M_Z \) and \( M_W \), and the weak-mixing angle, \( \theta \), to give the following\(^2\),

\[
F_2^\pm = F_2^\pm + k(-v_e \mp P_e a_e)F_2^\gamma Z + k^2(v_e^2 + a_e^2 \mp 2P_e v_e a_e)F_2^Z, \tag{2}
\]

\[
xF_3^\pm = k(-a_e \mp P_e v_e)xF_3^\gamma Z + k^2(2v_e a_e \pm P_e (v_e^2 + a_e^2))xF_3^Z, \tag{3}
\]

where \( k = \frac{1}{4\sin^2 \theta \cos^2 \theta \sqrt{Q^2 + M_Z^2}} \) and the vector and axial-vector coupling of the electron to the Z boson are \( v_e = -1/2 + 2\sin^2 \theta \) and \( a_e = -1/2 \) respectively. By taking the difference \( \bar{\sigma}^{-p} - \bar{\sigma}^+p \) one can extract \( xF_3 \) using unpolarised HERA-I data and net unpolarised data from HERA-II. As the polarisation dependence is removed, \( xF_3 \) can be written as,

\[
xF_3 = -a_e kxF_3^\gamma Z + 2v_e a_e k^2 xF_3^Z. \tag{4}
\]

Since the coupling \( v_e \) is small and \( k < 1 \), the interference term dominates \( xF_3 \). In leading order (LO) perturbative QCD the interference structure function can be explicitly written in terms of the valence quark distributions, \( u_v \) and \( d_v \),

\[
xF_3^\gamma Z = \frac{x}{3}(2u_v + d_v + \Delta), \tag{5}
\]

where \( \Delta = 2(u_{\text{sea}} - \bar{u} + c - \bar{c}) + (d_{\text{sea}} - \bar{d} + s - \bar{s}) \). Therefore \( xF_3^\gamma Z \) is determined by the valence quark distribution if the \( \Delta \) term is ignored, and is only weakly dependent on \( Q^2 \). To minimise statistical errors, the \( xF_3^\gamma Z \) measurements can be extrapolated in \( Q^2 \) and averaged in \( x \). Results from the ZEUS and H1 collaborations are shown in Fig. 1(a).
The Born level charged current (CC) $e^\pm p$ cross section with polarised leptons can be expressed at LO in QCD as\(^1\),

\[
\frac{d^2\sigma_{CC}(e^\pm p)}{dx dQ^2} = (1 - P_e) \frac{G_F^2}{2\pi} \left( \frac{M_W^2}{M_W^2 + Q^2} \right)^2 \left[ u + c \right] + (1 - y)^2(d + s),
\]

(6)

\[
\frac{d^2\sigma_{CC}(e^\pm p)}{dx dQ^2} = (1 + P_e) \frac{G_F^2}{2\pi} \left( \frac{M_W^2}{M_W^2 + Q^2} \right)^2 \left[ \bar{u} + \bar{c} \right] + (1 - y)^2(d + s),
\]

(7)

where $G_F$ is the Fermi coupling constant and $u, c, d, s$ are the respective quark densities. The flavour selecting nature of the CC interaction is apparent as $u$ quark content is revealed through $e^- p$ DIS, whereas $d$ quark constraints are possible through $e^+ p$ scattering. This can be illustrated in the $e^- p$ CC DIS reduced cross section measurements\(^4\) shown in Fig. 1(b), where the SM prediction describes the data well and is dominated by the $u$ quark density.

3 PDF Fits Using Only HERA Data and the Inclusion of Jet Data

The PDFs are usually determined in global fits using data from many different experiments. However, the high precision and wide kinematic coverage of existing HERA data allow precise extractions of the proton PDFs using only HERA data. The use of HERA data alone eliminates the uncertainty from heavy-target corrections and also avoids difficulties that can sometimes arise from combining data sets from several different experiments.

The high statistics HERA NC data is used to determine the low $x$ sea and gluon distributions while information on the valence quarks is provided by the higher-$Q^2$ NC and CC data. The gluon density contributes indirectly to the inclusive DIS cross sections, however it makes a direct contribution to the jet cross sections through boson-gluon fusion. The ZEUS collaboration has performed a combined NLO QCD fit (ZEUS-JETS PDF\(^5\)) to inclusive NC and CC DIS data as well as high precision jet data in DIS and $\gamma p$ scattering.

The ZEUS-JETS PDFs agree well with the previous ZEUS-S PDF global fits and are also compatible with the MRST\(^6\) and CTEQ\(^7\) PDFs. The shapes of the PDFs are not changed.
significantly by including jet data but the decrease in the uncertainty on the gluon distribution is significant, approximately halved, in the mid-$x$ region over the full $Q^2$ range.

4 Inclusion of New Data in PDFs and Future Projections from HERA

The PDF uncertainties from current global fits are, in general, limited by irreducible experimental systematics. In contrast, the fits to HERA data alone are largely limited by the statistical precision of existing measurements. Since 2003, HERA has delivered a substantial amount of luminosity with polarised lepton beams. Figure 2(a) shows a new PDF fit named ZEUS-pol$^8$ which includes HERA-II $e^-p$ NC and CC inclusive cross section data with a total integrated luminosity of 121.5 pb$^{-1}$. This leads to an improvement in PDF uncertainties at high $x$, especially for the $u$ valence quark.

As new HERA data is analysed, a significant impact on the gluon uncertainties could be made by future jet cross section measurements in kinematic regions optimised for PDF sensitivity. The effect on the PDF uncertainties using the entire HERA data set has been estimated in the HERA-II projection fit$^9$. A total integrated luminosity of 700 pb$^{-1}$ was assumed for the high $Q^2$ inclusive data, and 500 pb$^{-1}$ was assumed for the jet measurements with central values and systematic uncertainties taken from the published data in each case. A set of optimised jet cross sections were included for forward $\gamma p$ collisions assuming a luminosity of 500 pb$^{-1}$.

The increased statistical precision of the assumed amount of high $Q^2$ data gives a significant improvement in the valence quark uncertainties over the whole range of $x$. A significant improvement at high $x$ is seen for the sea quarks, however the low $x$ sea and low $x$ gluons are not significantly impacted as the data constraining this region tends to be at lower $Q^2$ and so already systematically limited. Much improvement is seen in the mid-to-high $x$ gluon which is constrained by jet data. Approximately half of the projected reduction in the gluon uncertainties is due to the inclusion of optimised jet cross sections.

Accurate proton PDFs are of great importance, especially for the LHC proton-proton collider which is planning to deliver high energy collisions in 2008. With HERA shutting down in July 2007, the projected improvements to the PDF uncertainties using solely HERA data will be particularly relevant to future physics at the LHC.

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QCD dynamics at low $x_{\text{Bj}}$ in $ep$ collisions at HERA

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Forward jet and multijet production has been measured at low Bjorken $x$ at HERA. The measured cross sections and correlations were compared to predictions from DGLAP-based fixed-order calculations. Further comparisons were made to DGLAP-based and CCFM-based leading-order Monte Carlo predictions, as well as to Colour-Dipole model predictions. For the majority of the phase space covered in the HERA kinematic region, fixed-order calculations describe the data well, while the leading-order models provide an inconsistent description of the data.

1 Introduction

Jet production in DIS is an ideal environment for investigating different approaches to parton dynamics at low Bjorken-$x$, $x_{\text{Bj}}$. An understanding of this regime is of particular relevance in view of the startup of the LHC, where many of the Standard Model processes such as the production of electroweak gauge bosons or the Higgs particle involve the collision of partons with a low fraction of the proton momentum.

In the usual collinear QCD factorisation approach, the cross sections are obtained as the convolution of perturbative matrix elements and parton densities evolved according to the DGLAP evolution equations. In these equations, all orders proportional to $\alpha_s \ln Q^2$ and terms with double logarithms $\ln Q^2 \cdot \ln 1/x$, where $x$ is the fraction of the proton momentum carried by a parton, which is equal to $x_{\text{Bj}}$ in the quark-parton model, are resummed. In the DGLAP approach, the parton participating in the hard scattering is the result of a partonic cascade ordered in transverse momentum, $p_T$. The partonic cascade starts from a low-$p_T$ and high-$x$ parton from the incoming proton and ends up, after consecutive branching, in the high-$p_T$ and low-$x$ parton entering in the hard scattering. At low $x_{\text{Bj}}$, where the phase space for parton emissions increases, terms proportional to $\alpha_s \ln 1/x$ may become large and spoil the accuracy of the DGLAP approach. In this phase-space region, a better description may come from the BFKL approach, which resums terms proportional to $\ln 1/x$, and the CCFM approach, which uses unintegrated gluon densities in an all-loop non-Sudakov resummation.

Parton evolution schemes at low $x_{\text{Bj}}$ were studied at HERA by measuring forward jet production and correlations in jet angles and transverse momentum. An excess of forward jets compared to DGLAP-based predictions and jets produced in the hard scatter that are not strongly correlated in transverse momentum may indicate the breakdown of DGLAP dynamics.

*On behalf of the H1 and ZEUS collaborations
2 Dijet Azimuthal Correlations

Dijet azimuthal correlations were investigated by the H1 Collaboration\(^1\) by measuring the cross-sections \(d^3\sigma/dx_Bd\Delta\phi^*\), where \(\Delta\phi^*\) is the azimuthal separation in the hadronic centre-of-mass (HCM) frame between the two selected jets closest to the scattered electron in pseudorapidity, \(\eta\). The measurements of \(\Delta\phi^*\) are reasonably well-described by NLOJET\(^2\) calculations at \(\mathcal{O}(\alpha_s^3)\), albeit within large theoretical uncertainties. To reduce the theoretical uncertainties, the measurements were normalised to the visible cross section for \(\Delta\phi^* < 170^\circ\). With a reduced theoretical uncertainty, the calculations are shown to predict a narrower \(\Delta\phi^*\) spectrum than is measured, especially at very low \(x_{Bj}\), as shown in Fig. 1. The measurements were also compared to predictions from two RAPGAP\(^3\) (DGLAP) samples, with one sample using only direct photons, the other using both direct and resolved photons; LEPTO\(^4\) (CDM); and CASCADE\(^5\) (CCFM). All models fail to describe \(\Delta\phi^*\) over the entire range in \(x_{Bj}\) covered.

3 Multijet Correlations

The sensitivity of parton evolution to the topology of the jet system was studied by the ZEUS collaboration\(^6\). Multi-differential cross sections as functions of the jet correlations in transverse momenta, azimuthal angles, and pseudorapidity have been measured for dijet and trijet production in the HCM frame. DGLAP-based calculations from NLOJET at \(\mathcal{O}(\alpha_s^2)\) and \(\mathcal{O}(\alpha_s^3)\) were compared to the measurements. The NLOJET calculations at \(\mathcal{O}(\alpha_s^2)\) do not describe the correlations in transverse momenta and azimuthal angle for dijet events; however with inclusion of higher-order terms, the NLOJET calculations at \(\mathcal{O}(\alpha_s^3)\) describe the dijet data over the entire range in \(x_{Bj}\) covered. The importance of higher-order terms at low \(x_{Bj}\) is seen especially when measuring the double-differential cross sections in \(Q^2\) and \(x_{Bj}\) for events with \(\Delta\phi_{HCM}^{j1,j2} < 120^\circ\), where \(\Delta\phi_{HCM}^{j1,j2}\) is the azimuthal separation of the two jets with the highest transverse energy. At low \(x_{Bj}\), the NLOJET calculations at \(\mathcal{O}(\alpha_s^3)\) are up to about one order of magnitude larger than the \(\mathcal{O}(\alpha_s^2)\) calculations and are consistent with the data, as seen presented in Fig. 2. The NLOJET calculations at \(\mathcal{O}(\alpha_s^3)\) also provide a reasonable description of the trijet measurements, with the description improving somewhat at higher \(x_{Bj}\).
4 Forward Jet Production

To examine the sensitivity of parton evolution to forward jet production, the ZEUS collaboration has studied jet production in an extended pseudorapidity range of $\eta_{LAB}^{jet} < 3.5$ by incorporating the Forward Plug Calorimeter (FPC) used during the HERAI running period. Measurements of cross sections as functions of $Q^2$, $x_B$, $E_T^{jet}$, and $\eta_{LAB}^{jet}$ are reasonably well-described by DGLAP-based calculations from DISENT, with large theoretical uncertainties at both low $x_B$ and high $\eta_{LAB}^{jet}$. Predictions from LEPTO (DGLAP); ARIADNE (CDM); and CASCADE, with two sets from the J2003 unintegrated gluon PDF used, were also compared to the measurements. Overall, ARIADNE provides the best description of the measured cross sections; LEPTO consistently underestimates the cross sections, and CASCADE fails to consistently reproduce the shapes of the distributions (see Fig. 3).

![Figure 3: ZEUS forward jets as a function of the kinematic variables $x_B$ and $Q^2$ compared to predictions from ARIADNE, LEPTO, and CASCADE.]

5 Trijet Production and Correlations

Trijet cross sections and correlations were measured by the H1 collaboration as a study of parton evolution at low $x_B$. Cross sections were measured as functions of $x_B$, jet pseudorapidity, scaled jet energies, and correlations in the jet angles $\theta'$ and $\psi'$. The variable $\theta'$ is defined as the angle between the proton beam and the jet with the highest transverse energy, while $\psi'$ is defined as the angle between the plane defined by the proton beam and the highest $E_T$ jet, and the plane defined by the two jets with the highest $E_T$. These measurements were made for three separate trijet samples: an inclusive trijet sample, and two trijet samples with one and two forward jets, respectively, with a forward jet having $\theta_{LAB}^{jet} < 20^\circ$ and $x_{jet} = E_{H,CMB}/E_{beam} > 0.035$. For the inclusive trijet sample, NLOJET calculations provide a reasonable description of the measured cross sections, but slightly underestimate the

![Figure 4: H1 trijet cross sections as a function of $x_B$ for events with two forward jets compared to NLOJET calculations at $O(\alpha_s^2)$ and $O(\alpha_s^3)$.]
measurements in the lowest bin of $x_{Bj}$. The agreement between the calculations and the measured cross section in $x_{Bj}$ is worse for the trijet sample containing two forward jets, with the most noticeable disagreement observed at lowest $x_{Bj}$ (see Fig. 4). The selection of two forward jets favors events with forward gluon emission unordered in transverse momentum, which the calculations at $\mathcal{O}(\alpha_s^3)$ do not predict entirely. Also seen in Fig. 4 is that the higher-order terms in the NLO JET calculations are important for forward jet emissions. The other cross sections for this sample are well-described by the calculations.

Predictions from DJANGOH (CDM) and RAPGAP LO MC models were also compared to the measured cross sections. The cross sections for the inclusive trijet sample are better described by CDM predictions, but both the CDM and RAPGAP predictions are inconsistent for the jet correlation angles $\theta'$ and $\psi'$; the RAPGAP predictions fail to describe the $\theta'$ distributions, and the CDM predictions fail to describe the $\psi'$ cross sections.

6 Summary

Parton dynamics at low $x_{Bj}$ ($10^{-4} < x_{Bj} < 10^{-2}$) have been investigated at HERA by the ZEUS and H1 collaborations. DGLAP-based NLO calculations describe the measured cross sections and jet correlations reasonably well for the most part when higher-order terms in the calculations are properly taken into account. The calculations fail to describe the trijet cross section in $x_{Bj}$ when the trijet sample contains two forward jets. Leading-order Monte Carlo models provide an inconsistent description of the measured cross sections. DGLAP-based LO MC models in general do not describe the cross sections; CDM models fail to describe $\Delta\phi^*$ and $\psi'$; CASCADE predictions are highly sensitive to the unintegrated gluon PDF used, and do not describe the data consistently.

References

FLUCTUATION PROPAGATOR AND HEAVY QUARK DIFFUSION

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The quark fluctuation propagator is evaluated.
It defines the diffusion coefficient in the vicinity of the phase transition
and the gradient term in the Ginzburg-Landau functional.

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1. Introduction

Recent experiments performed at RHIC have revealed an unexpectedly large heavy flavor suppression. This result may indicate that the heavy quark diffusion coefficient is anomalously small in the vicinity of the critical line of the QCD phase diagram in the $(\mu, T)$ plane. The diffusion coefficient enters also into the Ginzburg-Landau functional via the term proportional to the $< A_\mu^2 >$ condensate [1]. Comparison with the lattice calculations of $< A_\mu^2 >$ shows that due to some dynamical reasons the diffusion coefficient is much smaller than the value given by the simple Drude formula [1]. It is well known that the diffusion coefficient may become small in the fluctuation regime and turns zero at the Anderson localization edge.

The transport coefficients are expressed in terms of the time-dependent propagator. Below we draw the fluctuation quark propagator (FQP) in its simplest form as a preliminary step in the investigation of the heavy quark dynamics in the vicinity of the phase transition. Our derivation closely follows the guidelines of the condensed matter physics [2].

Denoting the FQP as $L(\tilde{p}, \omega)$ we may define it as

$$ L^{-1}(\tilde{p}, \omega) = -\frac{1}{g} + F(\tilde{p}, \omega), \quad (1) $$

$$ F(\tilde{p}, \omega) = \sum_k G(\vec{k}, k_4) G(\tilde{p} - \vec{k}, \omega - k_4), \quad (2) $$
Here $g$ is the coupling constant with the dimension $m^{-2}$, the sum over $k$ implies the momentum integration, Matsubara and discrete indices summation, $G(\tilde{k}, k_4)$ in the thermal Green's function which reads

$$G(\tilde{k}, k_4) = \frac{i}{\gamma \tilde{k} + \gamma_4 k_4 - \bar{\varepsilon} + i\mu \gamma_4},$$

(3)

with $k_4 = -\pi(2n_1 + 1)T, T = \beta^{-1}$. We shall compute $F(\vec{p}, \omega)$ in the approximation of long-wave fluctuations

$$F(\vec{p}, \omega) \simeq A(\omega) + B(\omega)\vec{p}^2,$$

(4)

First we compute the term $A(\omega)$. We have

$$tr\{G(\tilde{k}, k_4)G(-\tilde{k}, \omega - k_4)\} = 2 \left\{ \frac{1}{\tilde{k}_4^2 + (\varepsilon - \tilde{\mu})^2} + \frac{1}{\tilde{k}_4^2 + (\varepsilon + \tilde{\mu})^2} \right\},$$

(5)

where $tr$ is over the Lorentz indices, $\varepsilon^2 = \tilde{k}^2 + m^2, \tilde{k}_4 = k_4 - \omega/2, \tilde{\mu} = \mu - i\omega/2$. The second term in (5) corresponds to antiquarks. We shall omit it here only for brevity though as shown in [1] the interplay of the quark and antiquark modes may result in instability in the chiral limit. For the same reason we put $m = 0$. Performing the momentum integration around the Fermi surface we obtain

$$A(\omega) = \nu \sum_{n \geq 0} \frac{1}{n + \frac{1}{2} + \frac{\omega}{4\pi T}},$$

(6)

where $\nu = \frac{2\mu^2}{\pi^2}$ is the density of states at the Fermi surface for two quark flavors. To evaluate $B(\omega)$ we act by the operator $\left(\vec{p} \frac{\partial}{\partial k}\right)^2$ on the second Green's function in (2). The result reads

$$B = \frac{\nu}{48\pi^2 T^2} \sum_{n \geq 0} \frac{1}{n + \frac{1}{2} + \frac{\omega}{4\pi T}}^3,$$

(7)

The logarithmic divergence of the sum in (6) may be removed by the introduction of the critical temperature $T_c$. First we regularize (6) by introducing the high-frequency cut-off $n_{\text{max}} = \frac{\Lambda}{2\pi T}, \Lambda \gg \omega$. Then

$$A = \nu \left\{ \psi \left(\frac{1}{2} + \frac{\omega}{4\pi T} + \frac{\Lambda}{2\pi T}\right) - \psi \left(\frac{1}{2} + \frac{\Lambda}{2\pi T}\right) \right\},$$

(8)

where $\psi(z)$ is the logarithmic derivative of the $\Gamma$-function. Next we replace $\Lambda$ by the critical temperature $T_c$ using the relation
\[ t = \ln \frac{T}{T_c} = \frac{1}{\nu g} - \psi \left( \frac{1}{2} + \frac{\Lambda}{2\pi T} \right) - \psi \left( \frac{1}{2} \right). \] (9)

Now we can return to the underlying formula (1) for the FQP, collect all the pieces together and write

\[ L^{-1}(\bar{p}, \omega) = -\nu \left\{ t + \psi \left( \frac{1}{2} + \frac{\omega}{4\pi T} \right) - \psi \left( \frac{1}{2} \right) - \frac{\bar{p}^2}{96\pi^2 T^2} \psi'' \left( \frac{1}{2} + \frac{\omega}{4\pi T} \right) \right\}. \] (10)

Equation (10) is the basic one in condensed matter fluctuation theory [2]. We have shown that it can be almost literally retrieved within rather general approach to dense finite temperature quark matter. Fluctuations in quark matter are many orders of magnitude stronger than in ordinary and even in high temperature superconductors [1]. Quark matter formed in heavy ion collisions has a finite volume which also increases fluctuation effects. Close to the phase transition quark matter is a system with strong disorder [1] possibly revealing Anderson localization. The FQP is known to be an effective tool to study the properties of such systems. In particular the poles of \( L(\bar{p}, \omega) \) determine the dynamical diffusion coefficient. The detailed investigation of this problem is beyond the scope of the present quick paper.

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References

DIFFRACTION : RECENT RESULTS AND IMPLICATIONS FOR LHC

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With the knowledge of diffractive parton densities extracted from HERA data, we discuss the observation of exclusive events using the dijet mass fraction as measured by the CDF collaboration at the Tevatron. In particular the impact of the gluon density uncertainty is analysed. Some prospects are given for diffractive physics at the LHC.

1 Diffraction at HERA

Since years, the Pomeron remains a subject of many interrogations. Indeed, defined as the virtual colourless carrier of strong interactions, the nature of the Pomeron is still a real challenge. While in the perturbative regime of QCD it can be defined as a compound system of two gluons in the approximation of resumming the leading logs in energy, its non-perturbative structure is basically unknown. In the recent years, an interesting experimental investigation on “hard” diffractive processes led to a new insight into those problems. At the HERA accelerator, it has been discovered that a non negligible amount of $\gamma^* +$proton deep inelastic events can be produced with no visible breaking of the incident proton. There are various phenomenological interpretations of this phenomenon, but a very appealing one relies upon a partonic interpretation of the structure of the Pomeron. In fact, it is possible to nicely describe the diffractive cross-section data from HERA by a QCD DGLAP evolution of parton distributions in the Pomeron combined with flux factors describing phenomenologically the probability of finding a Pomeron state in the proton. Sets of quark and gluon distributions in the Pomeron following NLO $C^2$-evolution equations are obtained. The gluons are dominating the diffractive exchange and carry approximately 70 % of the momentum. The diffractive gluon density is presented in figure 1. At high $\beta$, where $\beta$ denotes the fraction of the particular parton in the pomeron, this density is not well constrained from the QCD fits. To quantify this uncertainty, we multiply the gluon distribution by the factor $(1-\beta^2)$ as shown in figure 1: we obtain the uncertainty on the parameter $\nu$, $\delta(\nu) = 0.5$, which corresponds to a large spread at large $\beta$ for the gluon density.

In the following, we investigate how this uncertainty influences the results on dijet mass fraction as measured at the Tevatron.

2 Diffraction at Tevatron and LHC

A schematic view of non diffractive, inclusive double pomeron exchange and exclusive diffractive events at the Tevatron or the LHC is displayed in figure 2. The upper left plot shows the "standard" non diffractive events where the Higgs boson, the dijet or diphotons are produced directly by a coupling to the proton associated with proton remnants. The bottom plot displays
the standard diffractive double Pomeron exchange (DPE) where the protons remain intact after interaction and the total available energy is used to produce the heavy object and the pomeron remnants. These events can be described using the diffractive gluon density measured at HERA and shown in figure 1. There may be a third class of processes displayed in the upper right figure, namely the exclusive diffractive production. Exclusive events allow a precise reconstruction of the mass and kinematical properties of the central object using the central detector or even more precisely using very forward detectors installed far downstream from the interaction point. The mass of the produced object can be computed using roman pot detectors and tagged protons, $M = \sqrt{s}\xi_1\xi_2$, where $\sqrt{s}$ is the energy of the reaction in the center of mass frame and $\xi_{1,2}$ represent the fractions of energy losses for both protons. We see immediately the advantage of these processes: we can benefit from the good roman pot resolution on $\xi_{1,2}$ to get a good resolution on mass. Therefore, it is possible to measure the mass and the kinematical properties of the produced object and use this information to increase the signal over background ratio by reducing the mass window of measurement.\(^3\)

If such exclusive processes exist in DPE, the most appealing is certainly the Higgs boson production through this channel at the LHC.\(^3\) It cannot be observed at the Tevatron due to the low production cross section, but one can use present measurements at the Tevatron to investigate any evidence for the existence of exclusive production in DPE.

3 Dijet mass fraction at the Tevatron

The CDF collaboration measured the so-called dijet mass fraction (DMF) in dijet events when the antiproton is tagged in the roman pot detectors and when there is a rapidity gap on the proton side to ensure that the event corresponds to a double pomeron exchange.\(^3\) The measured observable $R_{JJ}$ is defined as the ratio of the mass carried by the two jets divided by the total
Figure 2: Scheme of non-diffractive, inclusive double pomeron exchange and exclusive events at the Tevatron or LHC

diffractive mass. The DMF turns out to be a very appropriate observable for identifying the exclusive production, which would manifest itself as an excess of the events towards $R_{JJ} \sim 1$. Indeed, for exclusive events, the dijet mass is essentially equal to the mass of the central system because no pomeron remnant is present. Then, for exclusive events, the DMF is 1 at generator level and can be smeared out towards lower values taking into account the detector resolutions. The advantage of DMF is that one can focus on the shape of the distribution. The observation of exclusive events does not rely on the overall normalization which might be strongly dependent on the detector simulation and acceptance of the roman pot detector. Results are shown in figure 3 with Monte-Carlo expectations calculated using DPEMC$^5$. As we have seen in section 1, the uncertainty on the gluon density is large at large $\beta$, which directly reflects in different shapes for the DMF in the inclusive part. This is illustrated in figure 3 (left), where we show the impact of the parameter $x$, which quantifies the diffractive gluon density error, on the shape of the DMF. We observe that it is not sufficient to reproduce the behaviour of the DMF when $R_{JJ} \sim 1$. Indeed, we see a clear deficit of events towards high values of the DMF, where exclusive events are supposed to occur. In figure 3 (right), a specific model describing exclusive events$^6$ is added to the inclusive prediction and we obtain a good agreement between data and the sum of MC expectations$^6$. It is a first evidence that exclusive events have been observed at the Tevatron.

4 Dijet mass fraction at the LHC

The search for exclusive events at the LHC can be performed in the same channels as the ones used at the Tevatron. A direct precise measurement of the gluon density in the pomeron through the measurement of the diffractive dijet cross section at the LHC will be necessary to study in detail the exclusive events in the dijet channel and measure their cross section. This is why it is important to have the roman pots and the Silicon detectors installed during the
2009-2010 shutdown so that these measurements will allow to tune the models and the MC. On the other hand, it is also important to look for different methods to show the existence of exclusive events\(^3\). In addition, some other possibilities benefitting from the high luminosity of the LHC appear. One of the cleanest way to show the existence of exclusive events would be to measure the dilepton and diphoton cross section ratios as a function of the dilepton/diphoton mass. If exclusive events exist, this distribution should show a bump towards high values of the dilepton/diphoton mass since it is possible to produce exclusively diphotons but not dileptons at leading order.

Figure 3: Dijet mass fraction for jets \( p_T > 10 \, \text{GeV} \). The data are compared to inclusive model predictions including the uncertainty of the gluon density at high \( \beta \) (left) and to the sum of inclusive and exclusive predictions (right).

5 Conclusions

We have discussed a first evidence for the existence of exclusive events in double pomeron exchange at the Tevatron. If such events can be also observed at the LHC, it would be possible to produce a Higgs instead of a dijet system regarding the cross section values accessible at the LHC. The great benefit of exclusive events concerns the precise reconstruction of the mass of the central object, using roman pot detectors installed far downstream from the interaction point\(^3\). It gives the opportunity to work with a favorable signal/background ratio compared to standard Higgs searches with a mass below 150 GeV.

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INCLUSIVE AND EXCLUSIVE DIFFRACTION AT HERA

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Measurements of diffractive processes in ep collisions at HERA are presented. Sets of diffractive parton densities are obtained using inclusive diffractive sample and their universality is tested in diffractive dijet and charm production. Measurements of ρ0 and J/ψ in diffraction are also presented.

1 Inclusive Diffraction at HERA

At HERA ep collider the presence of processes of the type ep → eXp (Fig. 1) in deep-inelastic scattering (DIS) and low Björken x regime offers a controlled environment in which to study the QCD properties of diffraction. The diffractive structure of the proton is probed with a point-like photon. Diffractive processes are commonly discussed in terms of exchanges with net vacuum quantum numbers, though the exact nature of these exchanges is not well known. The detailed explanation of hard diffraction is a major challenge for the understanding of the strong interaction at high energies. In diffractive processes the proton exchanges a diffractive object (pomeron), carrying a small longitudinal momentum fraction xB, which couples to a photon of virtuality Q2, through a quark carrying the momentum fraction β = x/xB of the pomeron. The additional kinematic invariants are the 4-momentum transfer squared, t, at the proton vertex, the inelasticity, y, and the γp centre-of-mass energy, W. Such events can be selected experimentally either by requiring a large gap in rapidity between outgoing proton and the hadronic system.
$X$ or directly by measuring the outgoing proton using proton spectrometers. The main observable, the reduced diffractive cross section, $\sigma^D_r$, is given by 
\[
\frac{d^3\sigma^{p+p\rightarrow Xp}}{dx dy dz dt} = \frac{2\pi}{Q^2} Y \sigma^{D(4)}_r (x, Q^2, x_B, t),
\]
where $Y = (1 - y + y^2/2)$ and is related to the diffractive structure functions by $\sigma^D_r = F^D_2 \frac{1}{x} F^D_L$. A general theoretical framework is provided by the hard scattering QCD collinear factorisation theorem. This theorem states that diffractive structure functions can be factorised into universal diffractive parton distribution functions (DPDFs), $f^D_i$, of the proton (under the constraint of a leading final state proton with a particular 4-momentum) and the photon-quark cross section $\sigma^{\gamma q}, d\sigma^{p+p\rightarrow Xp} = \sum_i f^D_i \otimes \sigma^{\gamma q}$, where the sum runs over all contributing quark and anti-quark flavours. Furthermore, a proton vertex factorisation has been found to apply to a good approximation, $f^D_i(x, Q^2, x_B, t) = F_B f_i(\beta, Q^2)$, stating that the shape of the parton densities of the pomeron, $f_i$, are independent of the 4-momentum of the final state proton. The DPDFs obey DGLAP QCD evolution equations and can be extracted by applying QCD fits to diffractive structure function data.

2 Inclusive Measurements and Parton Densities

The H1 Collaboration has published recently the reduced cross section $\sigma^{D(3)}_r(x_B, \beta, Q^2)$ for the process $ep \rightarrow eXp$ under condition of proton dissociation system mass $M_Y < 1.6$ GeV and $|t| < 1$ GeV$^2$, for various fixed values of 0.0003 < $x_B$ < 0.03. The variation of $\sigma^{D(3)}_r(x_B)$ with $x_B$ and $t$ is parameterised by the pomeron flux factor $f_B = \alpha_B e^{B t} x_B^{1-2a x B}$ where the $\alpha_B(t) = \alpha_B(0) + B t$ is an effective pomeron trajectory with intercept $\alpha_B(0) = 1.118 \pm 0.008$ (exp) $+0.02/_{-0.03}$ (model) and $\alpha_B = 0.06 \pm 0.19$ GeV$^{-2}$. The $\alpha_B$ and the $t$-slope parameter, $B$, were obtained from fits to H1 FPS data (see below). The $\beta$ and $Q^2$ dependences of $\sigma^{D(3)}_r$ are interpreted in terms of DPDFs extracted from the NLO DGLAP QCD fit performed in the $\overline{MS}$ renormalisation scheme with massive charm and beauty quarks. The reduced diffractive cross section exhibits scaling violations with positive $\partial \sigma^D_r / \partial \ln Q^2$ up to high $\beta$. The fraction of the exchanged momentum carried by gluons integrated over the kinematic range is around 70%. At large $\beta$, the inclusive diffractive data are less sensitive to gluon density, confirmed by the presence of two alternative sets of DPDFs (Fig. 2) with different solutions for gluon density at high $\beta$ ("H1 2006 DPDF Fit A’ and ‘H1 2006 DPDF Fit B’).

Another set of DPDFs has been obtained by simultaneous fit to the inclusive diffractive data and the measurements of the diffractive dijet production, allowing for a sensitive determination of diffractive quark and gluon distributions in the range of 0.05 < $z$ < 0.9, $z$ being the fraction of the exchanged momentum carried by the parton entering the hard interaction. Quark and gluon distributions are parameterised at the starting scale $Q^2 = \mu^2_f = 2.5$ GeV$^2$ as $f(z, \mu^2_f) \equiv A_{i-z} B_{z(1-z)} C_{z}$ and are evolved to higher factorisation scales by a numerical solution of the NLO DGLAP evolution equations. For the gluon density at high $z$, the combined fit is compatible with the H1 2006 Fit B. For the gluon density at low $z$ and the singlet density, all three fits are in agreement.

The H1 Collaboration has also published recently the diffractive reduced cross section $\sigma^{D(4)}_r(\beta, Q^2, x_B, t)$ at $|t| = 0.25$ GeV$^2$ for the process $ep \rightarrow eXp$ where the outgoing proton
is measured by the Forward Proton Spectrometer (FPS). The data are compared with the results of H1 measurement using events selected on the basis of rapidity gap, which includes a contribution of proton dissociation to states with masses $M_Y < 1.6$ GeV. The ratio of the measured cross section to the cross section obtained using rapidity gap method is $1.23 \pm 0.03$ (stat) $\pm 0.16$ (syst), independently of $Q^2$, $\beta$ and $x_P$ within uncertainties. The $t$-dependence is parameterised by an exponential function such that $d\sigma/dt \propto e^{Bt}$. The resulting values of the slope parameter $B$ in the pomeron dominated range $x_P \leq 0.0094$ are close to 6 GeV$^{-2}$ and are independent of $x_P$ in this region.

3 Factorisation Properties of the DPDFs

Diffractive dijet and $D^*$ production measurements have been used to test the universality of DPDFs.

Measurement of diffractive process $ep \rightarrow eXY$, where system $X$ contains at least two jets has been measured in DIS regime ($4 < Q^2 < 80$ GeV$^2$) and in photoproduction ($Q^2 < 0.01$ GeV$^2$). The NLO prediction obtained using the 'H1 2006 Fit B' diffractive parton densities agree with the distributions in DIS regime within errors (Fig. 3a,b) and supports the validity of QCD hard scattering factorisation in diffractive DIS. A fit to the photoproduction data yields to a suppression factors of $0.47 \pm 0.16$ for the part of NLO calculation for which $x^2_{\gamma^*,PL} < 0.9$ and $0.53 \pm 0.14$ for $x^2_{\gamma^*,PL} > 0.9$, where $x^2_{\gamma^*,PL}$ is the fraction of the photon momentum entering the hard scattering and is reconstructed at the parton level from parton jets before hadronisation (Fig. 3c,d). The QCD hard scattering factorisation is therefore broken in photoproduction.

The diffractive open charm production in DIS and photoproduction regimes has been measured at HERA with two techniques used for the cross section measurements. In the first, the charm quark is tagged by the reconstruction of a $D^{*+}(2010)$ meson. In the second method, the displacement of tracks from the primary vertex characterised for charm decays is measured and used to separate the signal. This latter method is used for the measurement of the open charm contribution to the inclusive diffractive cross section in DIS. The measured $D^*$ cross section is compared to the results based on NLO prediction of diffractive extension of the programs HVQDIS in DIS and FMNR in photoproduction. The renormalisation and the factorisation scales are set to $\mu_r = \mu_f = \sqrt{Q^2 + 4m_c^2}$ in DIS and to $\mu_r = \mu_f = \sqrt{Q^2 + 4m_c^2}$ in photoproduction, with $m_c = 1.5$ GeV. Hadronisation corrections were evaluated using the LUND model. The prediction from the 'H1 2006 DPDF' fits describes data within large experimental errors and theoretical uncertainties, supporting the validity of QCD factorisation for open charm production in diffractive DIS and photoproduction. The charm contribution to the inclusive diffractive cross section at $Q^2 = 35$ GeV$^2$ and two different values of $x_P = 0.004$ and $x_P = 0.018$ is found to be $\approx 20\%$ on average, compatible with the charm fraction in inclusive DIS.

4 Vector Meson Final States

The diffractive photoproduction of $\rho$ mesons, $ep \rightarrow e\rho Y$, at large $|t|$ has been studied at HERA in the range of $Q^2 < 0.01$ GeV$^2$ and mass $M_Y < 5$ GeV. The $t$ dependence of the cross section
is measured in the range of $1.5 < |t| < 10.0$ GeV$^2$ and is described by a power law $d\sigma/dt \propto |t|^{-n}$ with $n = 4.26 \pm 0.06^{(\text{stat})} \pm 0.04^{(\text{syst})}$. Data indicate a violation of s-channel helicity conservation with contributions from both single and double helicity-flip being observed, in agreement with previous ZEUS measurement\textsuperscript{8}.

The proton-dissociative diffractive photoproduction of $J/\psi$ meson has been studied by the ZEUS Collaboration\textsuperscript{9}. The differential cross section $d\sigma(yp \rightarrow J/\psi Y)/dt$ has been measured in the range $1 < |t| < 20$ GeV$^2$, $50 < W < 150$ GeV and $z > 0.95$ and is reproduced by pQCD BFKL LL and nonL calculations\textsuperscript{10} with fixed strong coupling constant, while the predictions with running $\alpha_s$ are too steep. The data are well described by a model based on DGLAP evolution in the region of its validity, $|t| < M_{J/\psi}^2$. The cross section $\sigma^{p+p \rightarrow J/\psi} \gamma$ has been measured as a function of $W$ in four bins of $|t|$ and fitted to a dependence of $\sigma \propto W^5$. In the Regge formalism, $d\sigma/dt \propto W^4(\alpha_{p(t)}-1)$. A fit performed in four bins of $t$ yields to an intercept $\alpha_{p(t)} = 1.153 \pm 0.048^{(\text{stat})} \pm 0.039^{(\text{syst})}$ and a slope $\alpha'_{p(t)} = -0.020 \pm 0.014^{(\text{stat})} \pm 0.010^{(\text{syst})}$. The values are similar to H1 results\textsuperscript{11} and consistent with prediction of the BFKL pomeron\textsuperscript{12}.

A measurement of exclusive $p^0$ photoproduction\textsuperscript{13} in the kinematic range $20 < W < 90$ GeV and $|t| < 3$ GeV$^2$ has been performed by H1 Collaboration. The Pomeron trajectory was extracted from the $W$ dependence of the cross section, $\alpha_{p(t)} = 1.093 \pm 0.003^{(\text{stat})} \pm 0.007^{(\text{syst})}$ and $\alpha'_{p(t)} = 0.116 \pm 0.027^{(\text{stat})} \pm 0.030^{(\text{syst})}$. The result supports the measurement from ZEUS Collaboration\textsuperscript{14} that in the space-like region $t < 0$, the Pomeron trajectory has a significantly smaller slope than the value of $\alpha'_{p(t)} = 0.25$ GeV$^{-2}$, derived from hadron scattering data\textsuperscript{15}.

5 Conclusions

Diffractive deep inelastic ep scattering is described by a factorised cross section including a hard scatter matrix element and DPDFs. Diffractive photoproduction of dijets shows a breaking of QCD factorisation by a factor of $\approx 0.5$ in direct and resolved photon contributions. The charm contribution to inclusive diffractive cross section is found to be about 20%. The BFKL model describes the t-dependence of the $p^0$ and $J/\psi$ cross section at large $|t|$ values.

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Global analysis for determining fragmentation functions 
and their uncertainties in light hadrons

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Fragmentation functions are determined for the pion, kaon, and proton by analyzing charged-hadron production data in electron-positron annihilation. It is important that uncertainties of the determined fragmentation functions are estimated in this analysis. Analysis results indicate that gluon and light-quark functions have large uncertainties especially at small $Q^2$. We find that next-to-leading-order (NLO) uncertainties are significantly reduced in comparison with leading-order (LO) ones in the pion and kaon. The fragmentation functions are very different in various analysis groups. However, all the recent functions are roughly within the estimated uncertainties, which indicates that they are consistent with each other. We provide a code for calculating the fragmentation functions and their uncertainties at a given kinematical point of $z$ and $Q^2$ by a user.

Keywords: Fragmentation function, Quark, Gluon, QCD, Electron-positron annihilation

1 Introduction

A fragmentation function describes a hadronization process from a parent quark, antiquark, or gluon to a hadron. Hadron-production processes are often used for investigating important physics such as the origin of nucleon spin and properties of quark-hadron matters. Fragmentation functions are needed for describing such processes, so that precise functions should be obtained for discussing any physics outcome. Nevertheless, it is known that there are large differences in the parametrized fragmentation functions, for example, between the ones by Kniehl, Kramer, and Pötzler (KKP) and Kretzer. Recently updated functions by Albino, Kniehl, and Kramer (AKK) are also much different from these functions. This fact suggests that the fragmentation functions are not determined accurately; therefore, it is important to show reliable regions in discussing any hadron-production data.

Such error analyses have been investigated recently in the studies of unpolarized parton distribution functions (PDFs), polarized and nuclear PDFs. It is straightforward to apply the technique for the fragmentation functions. We determine the fragmentation functions and their uncertainties by analyzing the data for charged-hadron production in electron-positron annihilation, $e^+ + e^- \rightarrow h + X$. The analyses are done in leading order (LO) and next-to-leading order (NLO) of the running coupling constant $\alpha_s$. Because accurate SLD data in 2004 are
included in our analysis, whereas they are not used in KKP, AKK, and Kretzer’s analyses, we expect to have improvements. Therefore, important points of our analysis are:

- improvement due to addition of accurate SLD data,
- roles of NLO terms on the determination, namely on the uncertainties,
- comparison with other analysis results by considering the uncertainties.

Our analysis method is explained in section 2, results are explained in section 3, and they are summarized in section 4.

2 Analysis method

The fragmentation function is defined by the ratio of hadron-production cross section to the total hadronic cross section:

\[ F_h(z, Q^2) = \frac{1}{\sigma_{\text{tot}}} \frac{d\sigma(e^+ e^- \rightarrow hX)}{dz}, \]

where \( Q^2 \) is given by the center-of-mass energy squared \( (Q^2 = s) \), and \( z \) is defined by the ratio \( z = E_h/(\sqrt{s}/2) = 2E_h/\sqrt{Q^2} \) with the hadron energy \( E_h \). Since the fragmentation occurs from primary quarks, antiquarks, and gluons, the fragmentation function is expressed by the their sum:

\[ F_h(z, Q^2) = \sum_y \int_0^1 \frac{dy}{y} C_i(y, \alpha_s) D_i^h(z/y, Q^2). \]

Here, \( C_i(z, \alpha_s) \) is a coefficient function which is calculated in perturbative QCD, and \( D_i^h(z, Q^2) \) is the fragmentation function of the hadron \( h \) from a parton \( i \). The function \( D_i^h(z, Q^2) \) is associated with a non-perturbative aspect, and it cannot be theoretically calculated in a reliably way. It is the purpose of this work to obtain the optimum fragmentation functions for the pion, kaon, and proton by analyzing the experimental data for \( e^+ e^- \rightarrow h + X \).

In order to determine the functions from the data, we express them in terms of parameters at a fixed scale \( Q_0^2 \) (=1 GeV²):

\[ D_i^h(z, Q_0^2) = N_i^h z^{\alpha_i^h} (1 - z)^{\beta_i^h}, \]

where \( N_i^h, \alpha_i^h, \) and \( \beta_i^h \) are the parameters to be determined by a \( \chi^2 \) analysis of the data. Because there is a sum rule due to the energy conservation: \( \sum_h M_h^i = \sum_h \int_0^1 dz \int D_i^h(z, Q^2) = 1 \), it is more convenient to choose the parameter \( M_i^h \) instead of \( N_i^h \). They are related by \( N_i^h = M_i^h / B(\alpha_i^h + 2, \beta_i^h + 1) \), where \( B(\alpha_i^h + 2, \beta_i^h + 1) \) is the beta function. In general, a common function is assumed for favored functions and different ones are used for disfavored functions. The favored indicates a fragmentation from a quark or antiquark which exists in the hadron as a constituent in a simple quark model. The disfavored means a fragmentation from other quark or antiquark. The details of the formalism are explained in Ref. 5. The optimum parameters are determined by minimizing the total \( \chi^2 \) given by \( \chi^2 = \sum_j (F_j^{\text{data}} - F_j^{\text{theo}})^2 / (\sigma_{\text{data}}^2) \), where \( F_j^{\text{data}} \) and \( F_j^{\text{theo}} \) are experimental and theoretical fragmentation functions, respectively, and \( \sigma_{\text{data}}^2 \) is an experimental error. Uncertainties of the determined fragmentation functions are estimated by the Hessian method:

\[ [\delta D_i^h(z)]^2 = \Delta \chi^2 \sum_{j,k} \left( \frac{\partial D_i^h(z, \xi_j)}{\partial \xi_j} \right) H^{-1}_{jk} \left( \frac{\partial D_i^h(z, \xi_k)}{\partial \xi_k} \right), \]

where \( H_{jk} \) is the Hessian matrix, \( \xi_j \) is a parameter, \( \hat{\xi} \) indicates the optimum parameter set, and the \( \Delta \chi^2 \) value is chosen so that the error becomes the one-\( \sigma \) range in the multiparameter space. The detailed explanations for the uncertainties are found in Refs. 4 and 5.
3 Results

We explain analysis results. First, determined fragmentation functions are compared with charged-pion production data in Fig. 1. The curve indicates theoretical NLO results which are calculated by using determined parameters in the $\chi^2$ analysis, and the uncertainties are shown by the shaded band. The comparison suggests that the fit is successful in reproducing the data in four orders of magnitude.

Determined functions are shown at the initial scales ($Q^2=1$ GeV$^2$, $m_c^2$, and $m_b^2$) and also at an evolved scale $Q^2 = M_Z^2$ in Fig. 2. The LO and NLO functions and their uncertainties are shown. We notice that the uncertainties are generally large at small $Q^2$, especially in the LO. The gluon and light-quark functions have especially large uncertainties. However, it is interesting to note that the situation is much improved in the NLO because the uncertainties become significantly smaller. The uncertainty bands are smaller at large $Q^2$ ($= M_Z^2$). Since the fragmentation functions are used at small $Q^2$ ($\sim 1$ GeV$^2$), for example, in HERMES, RHIC-Spin, and RHIC heavy-ion experiments, one should be careful about the reliability of employed functions in one's analysis.

![Figure 1: NLO results are compared with charged-pion data.](image1)

![Figure 2: Determined fragmentation functions for $\pi^+$, $K^+$, and proton at $Q^2 = 1$ GeV$^2$, $m_c^2$, $m_b^2$, and $M_Z^2$. LO and NLO functions are shown with their uncertainties.](image2)

Next, the determined functions are compared with other analysis results for $(\pi^+ + \pi^-)/2$, $(K^+ + K^-)/2$, and $(p + \bar{p})/2$ in Fig. 3. Our parametrization is denoted HKNS (Hirai, Kumano, Nagai, Sudoh). The determined functions in NLO and their uncertainties are shown by the solid curves and shaded bands. They are compared with other functions by KKP, AKK, and Kretzer at $Q^2 = 2$, 10, and 100 GeV$^2$. As mentioned earlier, there are much differences between the analysis groups. For example, the gluon and $s$-quark functions have large variations in the pion. However, almost all the curves are roughly within the estimated uncertainty bands. It suggests that all the analyses should be consistent with each other and that accurate functions cannot be determined by the current $e^+e^-$ data. After our paper, there appeared another analysis by de Florian, Sassot, and Stratmann. Although there are some differences from our functions,
they are also within the uncertainty bands in Fig. 3.

The determined fragmentation functions can be calculated by using a code at our web site\textsuperscript{7} by supplying a kinematical condition for $z$ and $Q^2$ and a hadron species. It is noteworthy that the uncertainties can be also calculated by using the code.

4 Summary

The optimum fragmentation functions and their uncertainties have been obtained for the pion, kaon, and proton in both LO and NLO of $\alpha_S$ by the $\chi^2$ analyses of charged-hadron production data in electron-positron annihilation. It is the first analysis to show the uncertainties in the fragmentation functions. The uncertainties were estimated by the Hessian method. We found large uncertainties especially at small $Q^2$, so that they need to be taken into account for using the functions in the small $p_T$ regions of hadron-production measurements in lepton-proton, proton-proton, and heavy-ion reactions. We also found that the functions are determined more accurately in the NLO than the LO ones particularly in the pion by considering LO and NLO uncertainties. There are large differences between previous parametrizations of KKP, AKK, and Kretzer, but they are consistent with each other and with our results because they are within the uncertainty bands.

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7. A code for calculating the fragmentation functions and their uncertainties can be obtained from http://research.kek.jp/people/kumanos/fss.html.
Bottomonium and Charmonium at CLEO

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The bottomonium and charmonium systems have long proved to be a rich source of QCD physics. Recent CLEO contributions in three disparate areas are presented: (1) the study of quark and gluon hadronization using $\Upsilon$ decays; (2) the interpretation of heavy charmonium states, including non-$c\bar{c}$ candidates; and (3) the exploration of light quark physics using the decays of narrow charmonium states as a well-controlled source of light quark hadrons.

1 Introduction

The CLEO experiment at the Cornell Electron Storage Ring (CESR) is uniquely situated to make simultaneous contributions to both the bottomonium and charmonium systems in a clean $e^+e^-$ environment. Between 2000 and 2003 CLEO III ran with $e^+e^-$ center of mass energies in the $\Upsilon$ region. A subset of this period was spent below $B\bar{B}$ threshold, where $\approx 20M$, $\approx 10M$, and $\approx 5M$ decays of the $\Upsilon(1S)$, $\Upsilon(2S)$, and $\Upsilon(3S)$, respectively, were collected. In 2003, CESR lowered its energy to the charmonium region and the CLEO III detector was slightly modified to become CLEO-c. Since that time, there has been an energy scan from 3.97 to 4.26 GeV ($\approx 60\text{pb}^{-1}$), samples collected at $4170\text{GeV} (\approx 300\text{pb}^{-1}$, largely for $D_s$ physics) and the $\psi(3770)$ ($\approx 300\text{pb}^{-1}$, largely for $D$ physics), and a total of nearly 28M $\psi(2S)$ decays have been recorded, only 3M of which have been analyzed.

Three (of many) topics recently addressed by the CLEO collaboration will be discussed below. The reach is wide: from fragmentation in bottomonium decays, to the interpretation of heavy charmonium states, to the use of narrow charmonium states as a source of light quark hadrons.
2 Bottomonium and Fragmentation

The bottomonium system provides many opportunities to study the hadronization of quarks and gluons. The number of gluons involved in the decay of a bottomonium state can be controlled by the charge-conjugation eigenvalue of the initial state: the \( \Upsilon \) states decay through three gluons; the \( \chi_{bJ} \) states decay through two. In addition, the continuum – where \( e^+e^- \rightarrow q\bar{q} \) proceeds without going through a resonance – can be used as a source of quarks. Thus, particle production can be studied and compared in a number of different environments.

2.1 Quark and Gluon Fragmentation

In 1984, CLEO I first noticed an enhancement in baryon production in \( ggg \) (from \( \Upsilon(1S) \) decays) over \( q\bar{q} \) (from the \( e^+e^- \rightarrow q\bar{q} \) continuum), i.e., the number of baryons produced per \( \Upsilon(1S) \) decay was greater than the number produced per \( q\bar{q} \) continuum event. The interpretation of this phenomenon, however, was complicated by the fact that the \( ggg \) system consists of three partons (or three strings), while the \( q\bar{q} \) system only has two partons (or one string). A recent CLEO II analysis has confirmed these findings with greater precision and has extended the comparison beyond \( \Upsilon(1S) \) decays to the decays of the \( \Upsilon(2S) \) and \( \Upsilon(3S) \) states as well. Figure 1a shows new measurements of the enhancements of particle production in \( ggg \) over \( q\bar{q} \), where the “enhancement” of a particle species is defined as the ratio of the number of particles produced per event in \( \Upsilon \) decays to the number produced per event from the continuum. The ratio is binned in particle momentum and integrated. The MC predictions incorporate the JETSET 7.3 fragmentation model. In addition, the new analysis compares particle production in \( gg\gamma \) (radiative \( \Upsilon \) decays) and \( q\bar{q}\gamma \) (radiative continuum events). The comparison in this case is between systems both having two partons and one string. The energy of the radiated photon is used to monitor parton energies. Figure 1b shows the enhancements of \( gg\gamma \) over \( q\bar{q}\gamma \), where in this case the ratio is binned in the energy of the radiated photon and integrated. A few conclusions can be drawn from these studies: (1) baryon enhancements in \( ggg \) vs. \( q\bar{q}\gamma \) are somewhat smaller than in \( ggg \) vs. \( q\bar{q} \); (2) the number of partons is important, not just \( \sqrt{s} \); and (3) the JETSET 7.3 fragmentation model does not reproduce the data.
2.2 Anti-Deuteron Production

The production of (anti)deuterons in Υ decays provides another opportunity to study the hadronization of quarks and gluons. In this case, models predict that the gluons from the Υ decay first hadronize into independent (anti)protons and (anti)neutrons, which in turn "coalesce" into (anti)deuterons due to their proximity in phase space. CLEO has measured the production of anti-deuterons in Υ(1S) and Υ(2S) decays and has set limits on their production in Υ(4S) decays. The production of anti-deuterons is easier to measure experimentally than the production of deuterons since anti-deuterons are not produced in hadronic interactions with the detector and the small background makes them easy to spot using dE/dx in the drift chambers. The relative branching fraction of inclusive Υ(1S) → dX to Υ(1S) → ggg, gγ was found to be (3.36 ± 0.23 ± 0.25) × 10⁻⁵. For comparison, a 90% C.L. upper limit of anti-deuteron production in the continuum was set at 0.031 pb at √s = 10.5 GeV, which, given an hadronic cross-section of the continuum of around 3000 pb, results in less than 1 in 10⁶ q̅q̅ events producing an anti-deuteron. This is a factor of three less than what is seen in Υ(1S) decays.

3 Interpretation of Heavy Charmonium States

The past few years have seen something of a renaissance in charmonium spectroscopy with the discovery of the unexpected Υ(4260) and X(3872) states, among others. The Υ(4260) and X(3872), in particular, have been the source of much speculation due to their multiple sightings and the difficulties encountered in attempting to incorporate them into the conventional c̅c̅ spectrum. The contributions of CLEO to their interpretation will be discussed below. In addition, CLEO has recently made measurements pertaining to the charmonium character of the ψ(3770), which is more often used as a source of DD. While the ψ(3770) is well-known and has been assumed to be the expected 3D₁ state of charmonium, pinning down its properties contributes to our global understanding of the charmonium spectrum.

3.1 Υ(4260)

The Υ(4260) was first observed by BaBar⁶ decaying to π⁺π⁻J/ψ using e⁺e⁻ collisions with initial state radiation (ISR). This production mechanism requires the Υ(4260) have JPC = 1−−. However, there is no place for a vector with this mass in the conventional c̅c̅ spectrum. On one interpretation the Υ(4260) is a hybrid meson, a q̅q̅ pair exhibiting an explicit gluonic degree of
freedom. CLEO has made two recent contributions regarding the nature of the $Y(4260)$. First, an $e^+e^-$ energy scan was performed between 3.97 and 4.26 GeV. A rise in the production cross section was observed for both $\pi^+\pi^- J/\psi$ and $\pi^0\pi^0 J/\psi$ at 4.26 GeV in the ratio of roughly 2:1. This ratio suggests the $Y(4260)$ is an isoscalar. Second, CLEO (using CLEO III data in the $Y$ region) has confirmed the initial observation by BaBar in $\pi^+\pi^- J/\psi$ from ISR. This both confirms its existence and its $J^{PC} = 1^{--}$ nature. The measured mass and width, $4284^{+17}_{-16} \pm 4$ MeV/$c^2$ and $73^{+30}_{-25} \pm 5$ MeV/$c^2$, respectively, are also consistent with BaBar.

### 3.2 $X(3872)$

The $X(3872)$ was first observed by Belle in the reaction $B \to KX$, $X \to \pi^+\pi^- J/\psi$. It has subsequently been studied in several different channels by a variety of different experiments. From its decay and production patterns it likely has $J^{PC} = 1^{++}$. One of the most tantalizing properties of this state is that its mass is very close to $D^0 D^{*0}$ threshold, suggesting that it could be a $D^0 D^{*0}$ molecule or a four-quark state. Prior to the new measurement by CLEO, the binding energy of the $X(3872)$ ($M(D^0) + M(D^{*0}) - M(X(3872))$), assuming it to be a $D^0 D^{*0}$ bound state, was $-0.9 \pm 2.1$ MeV, where the error, perhaps surprisingly, was dominated by the mass of the $D^0$. CLEO improved this situation with a new precision $D^0$ mass measurement using the well-constrained decay $D^0 \to \phi K_S$ (Figure 2b) and found the mass to be $1864.847 \pm 0.150 \pm 0.095$ MeV/$c^2$. This results in a small positive binding energy $1 \sigma$ from zero: $+0.6 \pm 0.6$ MeV. This lends further credence to the molecular interpretation of the $X(3872)$.

### 3.3 $\psi(3770)$

The existence of the $\psi(3770)$ has been established for a long time. However, because it predominantly decays to $D\bar{D}$ its behavior as a state of charmonium has been relatively unexplored in comparison to its lighter partners. The electromagnetic transitions, $\psi(3770) \to \gamma \chi_{cJ}$, because they are straightforward to calculate, provide a natural place to study the charmonium nature of the $\psi(3770)$. CLEO has recently measured these transitions in two independent analyses.
In the first\textsuperscript{11}, the processes were measured by reconstructing the $\chi_{cJ}$ in their transitions to $\gamma J/\psi$ and then requiring the $J/\psi$ to decay to $e^+e^-$ or $\mu^+\mu^-$ (Figure 3a). In the second\textsuperscript{12}, the $\chi_{cJ}$ were reconstructed in several exclusive hadronic modes and then normalized to the process $\psi(2S) \rightarrow \gamma \chi_{cJ}$ using the same exclusive modes (Figure 3b). The first method favors the measurement of the transitions to $\chi_{c1}$ and $\chi_{c2}$ while the second method is more suited to the transition to $\chi_{c0}$. Combining the results of the two analyses, the partial widths of $\psi(3770) \rightarrow \gamma \chi_{cJ}$ were found to be $172 \pm 30$ keV for $J = 0$, $70 \pm 17$ keV for $J = 1$, and an upper limit of 21 keV at 90\% C.L. was set for $J = 2$. These measurements are consistent with relativistic calculations assuming the $\psi(3770)$ is the $3D_1$ state of charmonium.

4 Using Charmonium to Study Light Quarks

In addition to providing valuable information in its own right, the charmonium system can also serve as a well-controlled source of light quark states. While much effort has gone into the study of $\psi(2S)$ and $J/\psi$ decays (e.g. $J/\psi$ radiative decays to glueballs), the decays of the $\chi_{cJ}$ states are less familiar and hold complementary information. The $\chi_{cJ}$ states are produced proficiently through the reaction $\psi(2S) \rightarrow \gamma \chi_{cJ}$, with rates around 9\% for $J = 0$, 1, and 2, and can be reconstructed cleanly in many different decay modes in the CLEO detector.

As an exploratory study into the analysis of the resonance substructure of $\chi_{cJ}$ decays, CLEO has recently analyzed a series of three-body $\chi_{cJ}$ decays\textsuperscript{13} using approximately 3M $\psi(2S)$ events collected with the CLEO III and CLEO-c detectors. This anticipates the new sample of approximately 25M $\psi(2S)$ events. The decay modes analyzed include $\eta \pi^+\pi^-$, $K^+K^-\eta$, $K^+K^-\pi^0$, $p\bar{p}\pi^0$, $p\bar{p}\eta$, $\eta\pi^+\pi^-$, $K_SK^-\pi^+$, and $K^+\bar{p}\Lambda$. Branching fractions were measured to each of these final states, many for the first time. Figure 4 shows $\chi_{cJ}$ decays to three particularly well-populated final states. The $\chi_{c1}$ decays to $\eta \pi^+\pi^-$, $K^+K^-\pi^0$, and $K_SK^-\pi^+$ included sufficient statistics for a rudimentary Dalitz analysis. Figure 5 shows the results of a fit to the $\eta \pi^+\pi^-$ Dalitz plot using a crude non-interfering resonance model. Dominant contributions were found from $\alpha(980)\eta$, $f_2(1270)\eta$, and $\sigma\eta$ with fit fractions of $75.1 \pm 3.5 \pm 4.3\%$, $14.4 \pm 3.1 \pm 1.9\%$ and $10.5 \pm 2.4 \pm 1.2\%$, respectively. No evidence for new structures was found in either $\eta \pi^+\pi^-$ or the two $K\bar{K}\pi$ modes.

Studies analyzing $\chi_{cJ}$ substructure using the full CLEO sample of 28M $\psi(2S)$ decays are underway. One reaction that looks particularly promising is the decay $\chi_{c0} \rightarrow K\bar{K}\pi\pi$, which was shown to exhibit a rich substructure of $f$ and $K^*$ states in a recent BES analysis\textsuperscript{14}. 

Figure 4: Reconstructed $\chi_{cJ}$ states ($J = 0$, 1, and 2) from the reaction $\psi(2S) \rightarrow \gamma \chi_{cJ}$. From left to right, the $\chi_{cJ}$ states are reconstructed in the exclusive channels $p\bar{p}\pi^0$, $\eta \pi^+\pi^-$, and $K^+K^-\pi^0$ (from reference [13]).
Figure 5: The Dalitz plot and its projections from the decay $\chi_c1 \rightarrow \eta \pi^+\pi^-$. Overlaid is a fit to the resonance substructure using a crude non-interfering resonance model (from reference [13]).

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References

XLIInd RENCONTRES DE MORIOND

QCD and High Energy Hadronic Interactions

X - Theoretical Developments Session

Chairperson: Frank Krauss
THE DUAL MEISSNER EFFECT FROM GAUGE THEORIES WITH MINIMAL SUPERSYMMETRY

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In the last decade it became clear that methods and techniques based on supersymmetry (SUSY) provide deep insights in quantum chromodynamics and other supersymmetric and non-supersymmetric gauge theories at strong coupling. I review progress in understanding the color confining mechanism based on the Meissner effect in $\mathcal{N}=1$ SUSY theories which became possible after the recent discovery of non-Abelian flux tubes and confined monopoles.

At this conference a large number of talks were presented treating various particular processes where QCD plays a role. While these processes are of great practical importance and the results reported are impressive, they should not overshadow a Big Picture of which they all are just small bits and pieces. This Big Picture is in the making for 33 years, and still not yet complete. I will use this opportunity to briefly review recent progress. For a detailed account see [1], where one can find an exhaustive list of references.

Let us ask ourselves: what is the most remarkable feature of quantum chromodynamics and QCD-like theories? The fact that at the Lagrangian level one deals with quarks and gluons while experimentalists detect pions, protons, glueballs and other color singlet states — never quarks and gluons — is the single most salient feature of non-Abelian gauge theories at strong coupling. Color confinement makes colored degrees of freedom inseparable. In a bid to understand this phenomenon Nambu, 't Hooft and Mandelstam suggested in the mid-1970s (independently but practically simultaneously) a “non-Abelian dual Meissner effect.” At that time their suggestion was more of a dream than a physical scenario. According to their vision, “non-Abelian monopoles” condense in the vacuum resulting in formation of “non-Abelian chromoelectric flux tubes” between color charges, e.g. between a probe heavy quark and antiquark. Attempts to separate these probe quarks leads to stretching of the flux tubes, so that the energy of the system grows linearly with separation. That’s how linear confinement was visualized. However, at that time the notions of non-Abelian flux tubes and non-Abelian monopoles (let alone condensed
monopoles in non-Abelian gauge theories) were nonexistent. Nambu, ’t Hooft and Mandelstam operated with nonexistent objects.

One may ask where did these theorists get their inspiration from? There was one physical phenomenon known since long ago and well understood theoretically which yielded a rather analogous picture.

In 1933 Meissner discovered that the magnetic fields could not penetrate inside superconducting media. The expulsion of the magnetic fields by superconductors goes under the name of the Meissner effect. Twenty years later Abrikosov posed the question: “what happens if one immerses a magnetic charge and an anticharge in type-II superconductors (which in fact he discovered)?” One can visualize a magnetic charge as an endpoint of a very long and very thin solenoid. Let us refer to the $N$ endpoint of such a solenoid as a positive magnetic charge and the $S$ endpoint as a negative magnetic charge.

In the empty space the magnetic field will spread in the bulk, while the energy of the magnetic charge-anticharge configuration will obey the Coulomb $1/r^2$ law. The force between them will die off as $1/r^2$.

What changes if the magnetic charges are placed inside a large type-II superconductor?

Inside the superconductor the Cooper pairs condense, all electric charges are screened, while the photon acquires a mass. According to modern terminology, the electromagnetic U(1) gauge symmetry is Higgsed. The magnetic field cannot be screened in this way; in fact, the magnetic flux is conserved. At the same time the superconducting medium does not tolerate the magnetic field.

The clash of contradictory requirements is solved through a compromise. A thin tube is formed between the magnetic charge and anticharge immersed in the superconducting medium. Inside this tube superconductivity is ruined — which allows the magnetic field to spread from the charge to the anticharge through this tube. The tube’s transverse size is proportional to the inverse photon mass while its tension is proportional to the Cooper pair condensate. These tubes go under the name of Abrikosov vortices. In fact, for arbitrary magnetic fields he predicted lattices of such flux tubes. A dramatic (and, sometimes, tragic) history of this discovery is nicely described in Abrikosov’s Nobel Lecture.

Returning to the magnetic charges immersed in the type-II superconductor under consideration, one can see that increasing the distance between these charges (as long as they are inside the superconductor) does not lead to their decoupling — the magnetic flux tube becomes longer, leading to a linear growth of the energy of the system.

The Abrikosov vortex lattices were experimentally observed in the 1980s. This physical phenomenon inspired Nambu, ’t Hooft and Mandelstam’s ideas on non-Abelian confinement. Many people tried to quantify these ideas. The first breakthrough, instrumental in all current developments, came 30 years later, in the form of the Seiberg-Witten solution of $N = 2$ super-Yang-Mills [2]. This theory has eight supercharges which makes dynamics quite “rigid” and helps one to find the full analytic solution at low energies. The theory bears a resemblance to quantum chromodynamics, sharing common “family treats.” By and large, one can characterize it as QCD’s “second cousin.”

An important feature which distinguishes it from QCD is the adjoint scalar field whose vacuum expectation value triggers the spontaneous breaking of the gauge symmetry $SU(2) \rightarrow U(1)$. The ’t Hooft-Polyakov monopoles ensue. They are readily seen in the quasiclassical domain. Extended supersymmetry and holomorphy in certain parameters which is associated with it allows one to analytically continue in the domain where the monopoles become light — eventually massless — and then condense after a certain small deformation breaking $N = 2$ down to $N = 1$ is introduced. Next, at a much lower scale the (dual) U(1) gauge symmetry breaks, so that the theory is fully Higgsed. Electric flux tubes are formed.

This was the first ever demonstration of the dual Meissner effect in non-Abelian theory, a
celebrated analytic proof of linear confinement, which caused much excitement and euphoria in the community.

It took people three years to realize that the flux tubes in the Seiberg–Witten solution are not those we would like to have in QCD. Hanany, Strassler and Zaffaroni who analyzed [3] in 1997 the chromoelectric flux tubes in the Seiberg–Witten solution showed that these flux tubes are essentially Abelian (of the Abrikosov–Nielsen–Olesen type) so that the hadrons they would create would have nothing to do with those in QCD. The hadronic spectrum would be significantly richer. And, say, in the SU(3) case, three flux tubes in the Seiberg–Witten solution would not annihilate into nothing, as they should in QCD ... Ever since searches for genuinely non-Abelian flux tubes and non-Abelian monopoles continued, with a decisive breakthrough in 2003. By that time the program of finding field-theory analogs of all basic constructions of string/D-brane theory was in full swing. BPS domain walls, analogs of D branes, had been identified in supersymmetric Yang–Mills theory. It had been demonstrated that such walls support gauge fields localized on them. BPS saturated string-wall junctions had been constructed. And yet, non-Abelian flux tubes, the basic element of the non-Abelian Meissner effect, remained elusive.

They were first found in U(2) super-Yang–Mills theories with extended supersymmetry, $\mathcal{N} = 2$, and two matter hypermultiplets [5,6]. If one introduces a non-vanishing Fayet–Iliopoulos parameter $\xi$ the theory develops isolated quark vacua, in which the gauge symmetry is fully Higgsed, and all elementary excitations are massive. In the general case, two matter mass terms allowed by $\mathcal{N} = 2$ are unequal, $m_1 \neq m_2$. There are free parameters whose interplay determines dynamics of the theory: the Fayet–Iliopoulos parameter $\xi$, the mass difference $\Delta m$ and a dynamical scale parameter $\Lambda$, an analog of the QCD scale $\Lambda_{QCD}$. Extended supersymmetry guarantees that some crucial dependences are holomorphic, and there is no phase transition.

As various parameters vary, this theory evolves in a very graphic way, see Fig. 1. At $\xi = 0$ but $\Delta m \neq 0$ (and $\Delta m \gg \Lambda$) it presents a very clear-cut example of a model with the standard 't Hooft–Polyakov monopole [7]. The monopole is free to fly — the flux tubes are not yet formed.

Switching on $\xi \neq 0$ traps the magnetic fields inside the flux tubes, which are weak as long as $\xi \ll \Delta m$. The flux tubes change the shape of the monopole far away from its core, leaving the core essentially intact. Orientation of the chromomagnetic field inside the flux tube is essentially fixed. The flux tubes are Abelian.

With $|\Delta m|$ decreasing, fluctuations in the orientation of the chromomagnetic field inside the flux tubes grow. Simultaneously, the monopole which no longer resembles the 't Hooft–Polyakov monopole, is seen as the string junction [7]. It acquires a life of its own.

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*This program started from the discovery of the BPS domain walls in supersymmetric Yang–Mills theory [4]. Subsequent results mentioned in this paragraph are largely due to A. Yung and myself, for references see [1].
Finally, in the limit $\Delta m \to 0$ the transformation is complete. A global SU(2) symmetry restores in the bulk. Orientational moduli develop on the string worldsheet making it truly non-Abelian. The junctions of degenerate strings present what remains of the monopoles in this highly quantum regime. It is remarkable that, despite the fact we are deep inside the highly quantum regime, holomorphy allows one to exactly calculate the mass of these monopoles.

What remains to be done? The most recent investigations zero in on $\mathcal{N} = 1$ theories, which are much closer relatives of QCD than $\mathcal{N} = 2$. I have time to say just a few words on the so-called $\mathcal{M}$ model suggested recently [8] which at the moment seems very promising.

The unwanted feature of $\mathcal{N} = 2$ theory, making it less similar to QCD, is the presence of the adjoint scalar field. One can get rid of it making it heavy. To this end we must add a "meson" $\mathcal{M}$ superfield to $\mathcal{N} = 2$ theory, and an appropriately chosen superpotential. After the adjoint field is eliminated the theory has no 't Hooft–Polyakov monopoles in the quasiclassical limit. Nevertheless, a non-Abelian Meissner effect does take place: condensation of color charges (quarks) gives rise to confined monopoles. The very fact of their existence in $\mathcal{N} = 1$ supersymmetric QCD without adjoint scalars was not known previously. The analysis presented in Ref. [8] is analytic and is based on the fact that the $\mathcal{N} = 1$ theory under consideration can be obtained starting from $\mathcal{N} = 2$ SQCD in which the 't Hooft–Polyakov monopoles do exist, through a certain limiting procedure allowing one to track the status of these monopoles at various stages (analogous to the one described above and summarized in Fig. 1).

If a dual of $\mathcal{N} = 1$ SQCD with the gauge group $U(N)$ and $N_f = N$ quark flavors could be identified, the dualized $\mathcal{M}$ model would be equivalent to the demonstration of the non-Abelian dual Meissner effect.

The $\mathcal{M}$ model can be regarded as the first cousin of QCD since the adjoint fields typical of $\mathcal{N} = 2$ are eliminated in this theory. Even though supersymmetry is considerably weakened, the overall qualitative picture survives. This is probably one of the most important findings at the current stage.

And then, $\mathcal{N} = 0$ theories — sister theories of QCD — loom large ...

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References

Non-relativistic Quantum Mechanics versus Quantum Field Theories

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We briefly review the derivation of a non-relativistic quantum mechanics description of a weakly bound non-relativistic system from the underlying quantum field theory. We highlight the main techniques used.

In a first approximation, the dynamics of the Hydrogen atom can be described by the solution of the Schrödinger equation with a Coulomb potential. However, it is not always clear how to derive this equation from the more fundamental quantum field theory, QED, much less how to get corrections in a systematic way. A similar problem is faced in heavy quarkonium systems with very large heavy quark masses. In this situation the dynamics is mainly perturbative and one efficient solution to this problem comes from the use of effective field theories (EFTs) and in particular of pNRQCD\(^a\). This EFT takes full advantage of the hierarchy of scales that appear in the system (\(v\) is the velocity of the heavy quark in the center of mass frame and \(m\) is the heavy quark mass):

\[
m \gg mv \gg mv^2 \ldots
\]

and makes systematic and natural the connection of the Quantum Field Theory with the Schrödinger equation. Roughly speaking the EFT turns out to be something like:

\[
\left\{ \begin{array}{c}
\left( i \partial_0 - \frac{\mathbf{p}^2}{2m} - V_s^{(0)}(r) \right) \Phi(r) = 0 + \text{corrections to the potential} \\
\text{+ interaction with other low - energy degrees of freedom}
\end{array} \right\} \text{pNRQCD}
\]

where \(V_s^{(0)}(r) = -C_F\alpha_s/r\) in the perturbative case and \(\Phi(r)\) is the \(\bar{Q}-Q\) wave-function. This

\(^a\text{For a comprehensive review of pNRQCD see}$\^2.\)
EFT is relevant, at least, for the study of the ground state properties of the bottomonium system, non-relativistic sum rules and the production of $t$-$\bar{t}$ near threshold (for some recent applications see 6,7,8,9).

The key point in the construction of the EFT is to determine the kinematic situation we want to describe. This fixes the (energy of the) degrees of freedom that appear as physical states (and not only as loop fluctuations). In our case the degrees of freedom in pNRQCD are kept to have $E \sim mv^2$. In order to derive pNRQCD we sequentially integrate out the larger scales.

$$\begin{align*}
QCD & \quad \text{Integrating out the hard scale (m)} \\
\text{NRQCD}^{10} & \quad \text{Integrating out the soft scale (mv)} \\
pNRQCD & \quad E \sim mv^2
\end{align*}$$

In this paper, we would like to highlight the main techniques needed in order to perform efficiently high-precision perturbative computations in non-relativistic bound state systems. They can be summarized in four points:

1. Matching QCD to NRQCD: Relativistic Feynman diagrams
2. Matching NRQCD to pNRQCD (getting the potential): Non-Relativistic (HQET-like) Feynman diagrams
3. Observable: Quantum mechanics perturbation theory
4. Observable: Ultrasoft loops

The first two points explain the techniques needed to obtain pNRQCD from QCD, whereas the last two explain the kind of computations faced in the EFT when computing observables. All the computations can be performed in dimensional regularization and only one scale appears in each type of integral, which becomes homogeneous. This is a very strong simplification of the problem. In practice this is implemented in the following way:

**Point 1.** One analytically expands over the three-momentum and residual energy in the integrand before the integration is made in both the full and the effective theory11,12.

$$\begin{align*}
\text{QCD} & \quad \int d^4q f(q, m, |p|, E) = \int d^4q f(q, m, 0, 0) + \mathcal{O} \left( \frac{E}{m}, \frac{|p|}{m} \right) \sim C \left( \frac{\mu}{m} \right) \text{(tree level)} |_{NRQCD} \\
\text{NRQCD} & \quad \int d^4q f(q, |p|, E) = \int d^4q f(q, 0, 0) = 0 \text{!!}
\end{align*}$$

Therefore, the computation of loops in the effective theory just gives zero and one matches loops in QCD with only one scale (the mass) to tree level diagrams in NRQCD, which we schematically draw in the following figure:

---

\textsuperscript{b}It is also possible to study heavy quarkonium systems in the non-perturbative regime with pNRQCD profiting from the hierarchy of scales of Eq. (1), see\textsuperscript{8,9}. 
Point 2) works analogously\textsuperscript{13}. One expands in the scales that are left in the effective theory. We integrate out the scale $k$ (transfer momentum between the quark and antiquark). Again loops in the EFT are zero and only tree-level diagrams have to be computed in the EFT:

\[ \int d^4q f(q, k, |p|, E) = \int d^4q f(q, k, 0, 0) + O\left(\frac{E}{k}, \frac{|p|}{k}\right) \sim \delta h_a(\text{potential}) \] (3)

\[ \text{pNRQCD} \quad \int d^4q f(q, |p|, E) = \int d^4q f(q, 0, 0) = 0!! \] (4)

We illustrate the matching in the figure below. Formally the one-loop diagram is equal to the QCD diagram shown above. The difference is that it has to be computed with the HQET quark propagator \(1/(q^2 + i\epsilon)\) and the vertices are also different.

\[
\begin{align*}
\text{NRQCD} & \quad \sim \frac{\alpha}{k^2} \\
\text{pNRQCD} & \quad \sim \frac{\alpha^2}{m^2} (\ln k^2 + c) = \frac{1}{h_s^{(0)} - E} \quad \text{Vol} 
\end{align*}
\]

Once the Lagrangian of pNRQCD has been obtained one can compute observables. A key quantity in this respect is the Green function. In order to go beyond the leading order description of the bound state one has to compute corrections to the Green function ($\delta h_s$ schematically represents the corrections to the potential and $H_I$ the interaction with ultrasoft gluons):

\[ G_s(E) = \frac{1}{h_s^{(0)} + \delta h_s - H_I - E} = G_s^{(0)} + \delta G_s \quad G_s^{(0)}(E) = \frac{1}{h_s^{(0)} - E} . \]

These corrections can be organized as an expansion in $1/m$, $\alpha_s$ and the multipole expansion. Two type of integrals appear then, which correspond to points 3) and 4) above.
**Point 3.** For example, if we were interested in computing the spectrum at $O(m \alpha_s^2)$ (for QED $\text{se}^{d^4}$), one should consider the iteration of subleading potentials ($\delta h_s$) in the propagator:

$$
\delta G_s^{\text{pot.}} = \frac{\delta h_s}{h_s^{(0)} - E} + \frac{\delta h_s}{h_s^{(0)} - E} + \cdots \\
\sim \frac{1}{h_s^{(0)} - E} + \frac{1}{h_s^{(0)} - E} + \frac{1}{h_s^{(0)} - E} + \cdots
$$

At some point, these corrections produce divergences. For example, a correction of the type: $\delta(r)G_s^{(0)}(C_f\alpha_s/r)G_s^{(0)}\delta(r)$, would produce the following divergence:

$$
< r = 0 \frac{1}{E - p^2/m} C_f \frac{\alpha_s}{r} \frac{1}{E - p^2/m} | r = 0 > \\
\sim \int \frac{d^4p}{(2\pi)^4} \int \frac{d^4p^\prime}{(2\pi)^4} \frac{4\pi\alpha_s}{(p - p')^2 - mE} \frac{m}{(p - p')^2 - mE} \sim -C_f \frac{m^2\alpha_s}{16\pi} \left( \frac{1}{\epsilon} + 2 \ln \left( \frac{mE}{\mu_p} \right) + \cdots \right).
$$

Nevertheless, the existence of divergences in the effective theory is not a problem since they get absorbed in the potentials ($\delta h_s$). The same happens with ultrasoft gluons, **point 4**

$$
\delta G_s^{\text{us}} = \frac{1}{(E - V_0^{(0)} - p^2/m)} \sim G_c(E) \int \frac{d^4k}{(2\pi)^4} r \frac{k}{k + p^2/m + V_0^{(0)} - E} G_c(E)
$$

$$
\sim G_c(E) r \left( \frac{p^2/m + V_0^{(0)} - E}{\nu_\text{us}^2} \right)^3 \left( \frac{1}{\epsilon} + \gamma + \ln \left( \frac{p^2/m + V_0^{(0)} - E}{\nu_\text{us}^2} \right) + C \right) r G_c(E),
$$

which also produces divergences that get absorbed in $\delta h_s$. Overall, we get a consistent EFT.

**References**

Phase Diagram of QCD

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This is a brief summary of the contemporary understanding of the QCD phase diagram as a function of temperature $T$ and baryo-chemical potential $\mu_B$.

1 Introduction

Strong interactions are described by Quantum Chromodynamics (QCD) – a remarkable theory. QCD is a convincing practical example of the success of the quantum field theory. Asymptotic freedom allows QCD to be consistent down to arbitrary short distance scale, enabling us to define the theory completely in terms of the fundamental microscopic degrees of freedom – quarks and gluons. This fundamental definition is very simple, yet the theory describes a wide range of phenomena – from the mass spectrum of hadrons to deep-inelastic processes. As such, QCD should also possess well defined thermodynamic properties. The knowledge of QCD thermodynamics is essential for the understanding of such natural phenomena as compact stars and laboratory experiments involving relativistic heavy-ion collisions.

Full analytical treatment of QCD is very difficult. In certain limits, in particular, for large values of the external thermodynamic parameters temperature $T$ and/or baryo-chemical potential $\mu_B$, when thermodynamics is dominated by short-distance QCD dynamics, the theory can be studied analytically, due to the asymptotic freedom. But the most interesting experimental region of parameters $T$ and $\mu_B$ is that of order $\Lambda_{\text{QCD}}$ – the intrinsic QCD scale. This makes first principle lattice approaches, which do not rely on a small coupling expansion, an invaluable and the most powerful tool in studying QCD thermodynamics.

The full potential of lattice methods is close to being realized as far as the study of QCD at $\mu_B = 0$ is concerned. The status of thermodynamics of QCD at non-zero $\mu_B$ is different. The main impediment to lattice simulations is the notorious sign problem. No method devised so far is known, or expected, to converge to the correct physical result as the infinite volume limit is approached at fixed $\mu_B \neq 0$. However, since the most interesting structure of the QCD phase diagram (phase transitions and critical points) lie at nonzero $\mu_B$, any progress in this direction is especially valuable. Existing lattice methods generically rely on clever extrapolations from $\mu_B = 0$. These techniques yield interesting results in the regime of small, but already experimentally relevant $\mu_B$.

A contemporary view of the QCD phase diagram is shown in Fig. 1. It is a compilation of a body of results from model calculations, empirical nuclear physics, as well as first principle lattice QCD calculations and perturbative calculations in asymptotic regimes.

This report provides an overview of the structure of the QCD phase diagram based on available theoretical (lattice and model calculations) and phenomenological input.
2 The phase diagram

Thermodynamic properties of a system are most readily expressed in terms of a phase diagram in the space of thermodynamic parameters – in the case of QCD – as a $T\mu_B$ phase diagram. Each point on the diagram corresponds to a stable thermodynamic state, characterized by various thermodynamic functions, such as, e.g., pressure, baryon density, etc. (as well as kinetic coefficients, e.g., diffusion or viscosity coefficients, or other properties of various correlation functions).

Static thermodynamic quantities can be derived from the partition function – a Gibbs sum over eigenstates of QCD Hamiltonian, which can be alternatively expressed as a path integral in Euclidean space of the exponent of the QCD action.

2.1 Massless quark limit and chiral symmetry argument

In the chiral limit – the idealized limit when 2 lightest quarks, $u$ and $d$, are taken to be massless, the Lagrangian of QCD acquires chiral symmetry $SU(2)_L \times SU(2)_R$, corresponding to $SU(2)$ flavor rotations of $(u_L, d_L)$ and $(u_R, d_R)$ doublets independently. The ground state of QCD breaks the chiral symmetry spontaneously locking $SU(2)_L$ and $SU(2)_R$ rotations into a single vector-like $SU(2)_V$ (isospin) symmetry and generating 3 massless Goldstone pseudoscalar bosons – the pions. The breaking of the chiral symmetry is a non-perturbative phenomenon.

At sufficiently high temperature $T \gg \Lambda_{QCD}$, due to the asymptotic freedom of QCD, perturbation theory around the approximation of the gas of free quarks and gluons (quark-gluon plasma – QGP) should become applicable. In this regime chiral symmetry is not broken. Thus we must expect a transition from a broken chiral symmetry vacuum state to a chirally symmetric equilibrium state at some temperature $T_c \sim \Lambda_{QCD}$. The transition is akin to the Curie point in a ferromagnet – where the rotational $O(3)$ symmetry is restored by thermal fluctuations (chiral $O(4)$–$SU(2) \times SU(2)$ symmetry in QCD). Thermodynamic functions of QCD must be singular at the transition point – as always when the transition separates thermodynamic states of different global symmetry.

Thus, in the massless quark (chiral) limit, the region of broken chiral symmetry on the $T\mu_B$ phase diagram must be separated from the region of the restored symmetry by a closed boundary.
2.2 \( N_f = 2 \) chiral limit and tricritical point

For two massless quarks the transition can be either second or first order\(^1\). As lattice and model calculations show, both possibilities are realized depending on the value of the strange quark mass \( m_s \) and/or the baryo-chemical potential \( \mu_B \).

The point on the chiral phase transition line where the transition changes order is called tricritical point. The location of this point is one of the unknowns of the QCD phase diagram with 2 massless quarks. In fact, even the order of the transition at \( \mu_B = 0 \), which many older and recent studies suggest is of the second order is still being questioned (see review\(^2\)).

Neither can it be claimed reliably (model or assumption independently) that the transition, if it begins as a 2nd order at \( \mu_B = 0 \), changes to first order at larger \( \mu_B \). However, numerous model calculations show this is the case. Lattice calculations also support such a picture. Contemporary understanding of QCD at low \( T \) and large \( \mu_B \), recently reviewed in\(^3\), also point at a first order transition (at low-\( T \), high-\( \mu_B \)) from nuclear matter to color-superconducting quark matter phase.

2.3 Physical quark masses and crossover

When the up and down quark masses are set to their observed finite values, the diagram assumes the shape sketched in Fig. 1. The second order transition line (where there was one) is replaced by a crossover – the criticality needed for the second order chiral restoration transition requires tuning chiral symmetry breaking parameters (quark masses) to zero. In the absence of the exact chiral symmetry (broken by quark masses) the transition from low- to high-temperature phases of QCD need not proceed through a singularity. Lattice simulations do indeed show that the transition is a crossover for \( \mu_B = 0 \) (most recently and decisively Ref.\(^4\), see also Ref.\(^2\) for a review). Recent terminology for the QCD state near the crossover \( (T \sim (1-2)T_c) \) is strongly coupled quark-gluon plasma (sQGP).

Transport properties of sQGP have attracted considerable attention. For example, generally, the shear viscosity \( \eta \) is a decreasing function of the coupling strength. The dimensionless ratio of \( \eta/\sigma \) to the entropy density \( \sigma \) tends to infinity asymptotically far on either side of the crossover – in dilute hadron gas \( (T \rightarrow 0) \) and in asymptotically free QGP \( (T \rightarrow \infty) \). Near the crossover \( \eta/\sigma \) should thus be expected to reach a minimum\(^5\). The viscosity can be indirectly determined in heavy ion collisions by comparing hydrodynamic calculations to experimental data. Such comparison\(^6\) indeed indicates that the viscosity (per entropy density) of this “crossover liquid” is relatively small, and plausibly is saturating the lower bound conjectured in Ref.\(^7\).

2.4 Physical quark masses and the critical point

The first order transition line is now ending at a point known as the QCD critical point or end point\(^8\). The end point of a first order line is a critical point of the second order. This is by far the most common critical phenomena in condensed matter physics. Most liquids possess such a singularity, including water. The line which we know as the water boiling transition ends at pressure \( p = 218 \text{ atm} \) and \( T = 374^\circ \text{C} \) in a critical point.

Beside the critical point, the phase diagram of QCD in Fig. 1 has other similarities with the phase diagram of water. A number of ordered quark matter phases must exist in the low-\( T \), high-\( \mu_B \) region, which are akin to many (more than 10) confirmed phases of ice. For asymptotically large \( \mu_B \), QCD with 3 quark flavors must be in color-flavor locked (CFL) state\(^8\).

\footnote{This fact is technically easier to establish than the order of the transition in the chiral limit – taking the chiral limit is an added difficulty.}

\footnote{The QCD critical point is sometimes also referred to as chiral critical point which sets it apart from another known (nuclear) critical point, the end-point of the transition separating nuclear liquid and gas phases (see Fig. 1). This point occurs at much lower temperatures \( O(10 \text{MeV}) \) set by the scale of the nuclear binding energies.}
3 Locating the critical point: the sign problem

The critical point is a well-defined singularity on the phase diagram, and it appears as an attractive theoretical, as well as experimental, target to shoot at. Theoretically, finding the coordinates \((T_c, \mu_B)\) of the critical point is a straightforwardly defined task. We need to calculate the partition function of QCD and find the singularity corresponding to the end of the first order transition line. But it is easier said than done.

Of course, calculating such a path integral analytically is beyond present reach. Numerical lattice Monte Carlo simulations is an obvious tool to choose for this task. At zero \(\mu_B\) Monte Carlo method allows us to determine the equation of state of QCD as a function of \(T\) (and show that the transition is a crossover). However, at finite \(\mu_B\) the Nature guards its secrets better.

The notorious sign problem has been known to lattice Monte Carlo experts since the early days of this field. Calculating the partition function using Monte Carlo method hinges on the fact that the exponent of the Euclidean action \(S_E\) is a positive-definite function of its variables (values of the fields on the lattice). This allows one to limit calculation to a relatively small set of field configurations randomly picked with probability proportional to the value of \(\exp(-S_E)\).

In QCD with \(\mu_B \neq 0\) the Monte Carlo action \(S_{\text{MC}}\) (playing the role of \(S_E\)) is complex. With \(S_{\text{MC}}\) complex, how does one pick configurations? A number of ways to circumvent the problem have been tried. For example, using the modulus of \(\exp(-S_{\text{MC}})\) as a probability measure, or the value of \(\exp(-S_{\text{MC}})\) at zero \(\mu_B\), when it is still positive. Unfortunately, none of the methods can be expected to converge to correct results with the increasing lattice volume \(V\), unless this limit is accompanied by an exponential \(\exp(\text{const}.\cdot V)\) increase of the number of configurations, rendering Monte Carlo techniques useless.

In the absence of a reliable first-principle approach model calculations have been the main source of knowledge about the QCD phase diagram\(^{9,10,11,12,13,14,15,16,17,18}\). This situation has began to change recently.

4 Lattice approaches to finding the critical point

This section is devoted to brief (and necessarily incomplete) descriptions of currently developed lattice methods for reaching out into the \(T\mu_B\) plane. For a more comprehensive description of these methods, as well as other methods not discussed here, the reader may consult the most up-to-date reviews\(^{19,20}\) as well as an earlier review by Philipsen\(^{21}\), which also contain further references to original papers.

4.1 Reweighting

The first lattice prediction for the location of the critical point was reported by Fodor and Katz in Ref.\(^{22}\). The assumption is that, although the problem becomes exponentially difficult as \(V \rightarrow \infty\), in practice, one can get a sensible approximation at finite \(V\). In addition, simulations at finite \(T\) might suffer lesser overlap problem because of large thermal fluctuations\(^{23}\). One can hope that if the critical point is at a small value of \(\mu_B\), the volume \(V\) may not need to be too large to achieve a reasonable accuracy.

4.2 Imaginary \(\mu_B\) and \(N_l = 3\)

By the universality argument of Ref.\(^{1}\), the finite temperature transition is 1st order for \(m_u = m_d = m_s = 0\). By continuity, it must remain 1st order in a finite domain of the \(m_s m_{ud}\) plane (taking \(m_u = m_d \equiv m_{ud}\) surrounding the origin – the plot of this domain is known as Columbia plot\(^{24,25}\). For physical quark masses and \(\mu_B = 0\) the temperature driven transition is a crossover,
which means that the physical point is outside of the 1st order domain in the $m_s m_{ud}$ plot. Reducing quark masses should pull the point into the 1st order domain.

What happens to the critical point when $(m_s, m_{ud})$ is in the 1st order domain? It is still a singularity of the partition function as a function of $\mu_B$, but it moves out into the complex $\mu_B$ plane. More precisely, it moves onto imaginary $\mu_B$ axis. This remarkable fact allows one to look at (the complex descendant of) the critical point in a direct Monte Carlo simulation – since there is no sign problem for imaginary $\mu_B$. This observation is at the core of the method developed by de Forcrand and Philipsen\textsuperscript{25,26}.

4.3 Taylor expansion

Taylor expansion in $\mu_B$ is another method to circumvent the sign problem. Derivatives of pressure (or other thermodynamic quantities) are calculated at $\mu_B = 0$ and assembled into a Taylor series expansion to obtain dependence of that quantity on $\mu_B$\textsuperscript{27,28,29}.

At fixed temperature, the convergence radius of the Taylor expansion in $\mu_B$ is limited by the nearest singularity in the complex plane of $\mu_B$. Assuming that at the temperature $T_E$, at which the critical point $(T_E, \mu_E)$ occurs on the phase diagram, this critical point is the nearest singularity to $\mu_B = 0$, one could use Taylor expansion to determine $\mu_E$\textsuperscript{27,28,29,30}, if $T_E$ is known.

Assuming that the radius of convergence $\mu_E$ can be approximated using the first few terms of the Taylor expansion, the main remaining problem is to determine the value of $T_E$ i.e., to identify at which value of $T$ the complex singularity reaches the real axis in the $\mu_B$ plane. This question has been addressed using universality arguments, as well as an example random matrix calculation in Ref\textsuperscript{30}.

5 Scanning QCD phase diagram in heavy ion collisions

Even though the exact location of the critical point is not known to us yet, the available theoretical estimates suggest that the point is within the region of the phase diagram probed by the heavy-ion collision experiments. This raises the possibility to discover this point in such experiments\textsuperscript{31}.

It is known empirically that with increasing collision energy, $\sqrt{s}$, the resulting fireballs tend to freeze out at decreasing values of $\mu_B$, i.e., decreasing baryon-antibaryon asymmetry. This is easy to understand, since the amount of generated entropy (heat) grows with $\sqrt{s}$ while the net baryon number is limited by that number in the initial nuclei.

The information about the location of the freezeout point for given experimental conditions is obtained by measuring the ratios of particle yields (e.g., baryons or antibaryons to pions), and fitting to a statistical model with $T$ and $\mu_B$ as parameters\textsuperscript{32}.

As with any critical point, measurement of fluctuations can be used to determine when the system is in the vicinity of the critical point. By measuring variables sensitive to the proximity of the critical point as a function of monotonically increasing $\sqrt{s}$ of the collision, and observing non-monotonic dependence, one discovers the critical point\textsuperscript{31}. The values of $T \mu_B$ corresponding to the freezeout at such a value of $\sqrt{s}$ give the coordinates of the critical point.

Acknowledgments

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References

LHC PREDICTIONS FOR TOTAL CROSS-SECTIONS FROM THE EIKONAL MINI-JET MODEL AND THE FROISSART BOUND

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We present results for total cross-section at LHC using an eikonal mini-jet model, where the rise is driven by QCD jets, and soft gluon effects are responsible for taming the rise to satisfy the Froissart bound and into the observed smooth increase.

1 Introduction

We shall first describe our model and then apply it to obtain a set of predictions at LHC both for total cross-section as well as for the survival probability of Large Rapidity Gaps. The connection between this model and the limits imposed by the Froissart bound are discussed.

1.1 The Bloch-Nordsieck (BN) model

The model we use\textsuperscript{1} is based on

1. hard component of scattering responsible for the rise of the total cross-section\textsuperscript{2,3}

2. soft gluon emission from scattering particles which softens the rise\textsuperscript{4}

3. eikonal transformation which implies multiple scattering and requires impact parameter distributions inside scattering particles and basic scattering cross-sections\textsuperscript{5}

With this model one obtains the total cross-section as shown in Fig. 1. According to our model, soft gluon emission is responsible for the initial decrease in pp, as well as for the transformation of the sharp rise due to the increase in gluon-gluon interactions into a smooth behavior. The
Figure 1: Total cross section data for pp and $p\bar{p}$ scattering compared with predictions from BN model

Figure 2: The minijet cross-section for different partons densities

The hard component of scattering responsible for the rise of the total cross-section is shown in Fig. 2 and it is obtained from the usual LO expression for the QCD jet cross-section

$$\sigma_{\text{hard}} \equiv \sigma_{\text{jet}}^A (s) =$$

$$\int_{\frac{\sqrt{s}}{2}}^{\sqrt{s}} \int_{0}^{1} dx_1 \int_{0}^{1} dx_2 \sum_{i,j,k,l} f_{i/A}(x_1) f_{j/B}(x_2) \frac{d\sigma^{kl}_{ij}(s)}{dp_t}$$

where $f_{i/A}(x_1)$ are PDF’s obtained from current parametrizations such as GRV, MRST, or CTEQ.

1.2 Producing the curves

We use the following eikonal expression for the total inelastic cross-section:

$$\sigma_{\text{inel}} = \int d^2\vec{b}[1 - e^{-n(b,s)}]$$

(1)

where $n(b,s)$ is the average number of inelastic collisions at an impact parameter $b$. Neglecting the real part of the eikonal, we then calculate the total cross-section as

$$\sigma_{\text{total}} = 2 \int d^2\vec{b}[1 - e^{-n(b,s)/2}]$$

(2)

In our model $n(b,s)$ is split as

$$n(b,s) = n_{\text{soft}}(b,s) + n_{\text{hard}}(b,s)$$

(3)

where we assume the following generic form

$$n_{\text{soft/hard}}(b,s) = A_{BN}^{\text{soft/hard}}(b,s) \sigma_{\text{soft/hard}}(s)$$

(4)

with

$$A_{BN}(b,s) = N \int \frac{d^2 K}{(2\pi)^2} e^{-iK \cdot b} \frac{d^2 P(K_\perp)}{d^2 K_\perp}.$$
Here $N$ is a normalization factor such that \( \int d^2 \vec{b} A(b) = 1 \) and

\[
\frac{d^2 P(K_\perp)}{d^2 K_\perp} = \frac{1}{(2\pi)^2} \int d^2 b e^{i K_\perp \cdot b} h(b, q_{\text{max}})
\]

is the transverse momentum distribution of initial state soft gluons emitted in the parton-parton collisions with

\[
h(b, q_{\text{max}}) = \int_0^{q_{\text{max}}} d^3 \vec{n}(k)[1 - e^{-\Delta k T}]
\]

\[
\approx \int_0^{q_{\text{max}}} \frac{8\alpha_s(k_T^2)}{3\pi} \frac{dq_{\text{max}}}{k_T} \log \frac{2q_{\text{max}}}{k_T} [1 - e^{-\Delta k T}]
\]

$q_{\text{max}}$ is the maximum transverse momentum allowed by kinematics to single soft gluon emission in a given hard collision, averaged over the parton densities and is obtained from the expression

\[
q_{\text{max}}(s) = \frac{\sqrt{s}}{2} \sum_{i,j} \int \frac{d x_1}{x_1} f_i(x_1) \int \frac{d x_2}{x_2} f_j(x_2) \sqrt{x_1} x_2 f_{i_{\text{min}}}^1 \frac{d z}{z} (1 - z)
\]

The steps we follow to compare the model with data are then the following:

1. choose the parameters for the hard scattering part, namely
   (i) parton densities (PDF), $p_{\text{min}}$ and $\Lambda_{\text{QCD}}$ for the chosen PDF set in Equation for $\sigma_{\text{jet}}$
   (ii) model for $\alpha_s$ in the infrared region and relevant parameters in the equation for soft emission

2. calculate $q_{\text{max}}(s, p_{\text{min}})$ for the given densities and $p_{\text{min}}$ using Equation for $q_{\text{max}}$

3. calculate $n_{\text{hard}}(b, s) = A_{B/N}^{\text{hard}}(b, s) \sigma_{\text{jet}}(s, p_{\text{min}})$

4. choose the parameters for the low energy part, namely
   (i) the constant low energy cross-section $\sigma_0$
   (ii) values for $q_{\text{max}}^{\text{soft}}$

5. calculate $n_{\text{soft}}(b, s) = A_{B/N}^{\text{soft}}(b, s) \sigma_0(1 + \epsilon \frac{\ln b}{\sqrt{s}})$ with $\epsilon = 0, 1$ depending upon the process being $pp$ or $pp$

6. calculate $n(b, s)$ and thus $\sigma_{\text{tot}}$

7. choose the parameter set which gives the best description of the total cross-section up to the Tevatron data

1.3 Survival probability for Large Rapidity Gaps and total cross-section

Once the $b$-distribution is available, one can use it to calculate the probability of not having an inelastic collision, a quantity easily available in the eikonal models since it is given by

\[
P_{\text{inel}} = e^{-n(b, s)}
\]

It can be used as the survival probability of large rapidity gaps for collisions at a given $b$

\[
< |S|^2 > = \int (d^2 b) A(b, q_{\text{max}}) |S(b)|^2
\]

In Fig. 3 and Fig. 4 we show the predictions of this model for two quantities of interest at LHC, the Survival probability and the total cross-section, and compare them other models.
2 About the Froissart bound and QCD minijets

The QCD jet cross-section for a fixed minimum value of the outgoing quark transverse momentum $p_T$ increases very rapidly as the c.m. energy of the colliding hadrons increases. For all densities we find

$$\sigma_{jet}^{PDF}(s, p_{\text{trans}}) \approx \epsilon,$$

with $\epsilon \approx 0.4$ for GRV and GRV98 which are more singular and $\epsilon \approx 0.3$ for CTEQ and MRST, which are less singular. Embedding the mini-jet cross-section in an eikonal reduces the growth of the minijet cross-sections since the eikonal representation produces a finite size to the scattering hadrons. The finite size may be introduced as follows:

- An energy independent b-distribution as in the form factor representation,

$$A(b) \sim e^{-b \times \text{constant}} \quad \text{at very large } b.$$ (9)

But the above has been found not enough to tame the rise because the growth of $\sigma_{jet}$ is too strong,

- Energy dependent distribution as in our model where satisfaction of the Froissart bound implies that the finite range of the interaction is restored through soft gluon emission

A variation of the parameters and the employment of different PDF with soft gluon emission along with QCD minijets does lead to an energy dependent size of the cross-section grows at most like $\log^2{s}$.

An example of how the Froissart bound is implemented in a mini-jet model with soft gluon emission is illustrated below.

$$h(b, q_{\text{max}}) \sim b^2 \varphi(b) \quad \text{at } b \approx \sqrt{s}$$ (10)

then $[1 - e^{-\varphi(b)}]$ behaves as a Fermi function and

$$\sigma_{tot} = 2 \int (d^2b) \varphi(b_0 - b) = 2\pi b_0^2$$ (11)
where $b_0$ is that value of $b$ for which $[1 - e^{-n(b_0,s)/2}] = 1/2$, or, equivalently

$$e^{-n(b_0,s)/2} = 1/2 \quad \text{or} \quad n(b_0,s) = 2 \log 2$$

(12)

$$n_{hard}^{PDF}(b,s) \approx A_{hard}^{PDF}(b,s) n_{jet}^{PDF}(s, p_{min})$$

(13)

$$n_{hard}^{PDF}(b,s) \approx c(s) e^{-b^2 c(s)} \sigma_{jet}^{PDF}(s, p_{min})$$

(14)

$$\sigma_{total} \approx 2 \pi b_0^2 = \frac{\pi}{c(s)} \log \frac{c(s) \sigma_{jet}^{PDF}(s, p_{min})}{2 \ln 2}$$

(15)

at very large $s$.

- $\sigma_{jet}^{PDF}(s, p_{min}) \sim s^c$ as expected from an infinite range theory like QCD
- but if $c(s) \sim constant \ or \ \uparrow \sqrt{s}$

We finally obtain

$$\sigma_{total} \leq \log s \ as \sqrt{s} \uparrow$$

(16)

No violation of the Froissart bound implies that the finite range of the interaction has indeed been restored through the soft gluon emission.

XI - Heavy Ion Session

Chairperson: M. Stephanov and Pamela Ferrari
THE IMPORTANCE OF THE INITIAL GEOMETRY
IN HEAVY ION COLLISIONS

R.S. HOLLIS\textsuperscript{6} for the PHOBOS Collaboration
B.Alver\textsuperscript{4}, B.B.Baack\textsuperscript{1}, M.D.Baker\textsuperscript{2}, M.Ballintijn\textsuperscript{4}, D.S.Barton\textsuperscript{2}, R.R.Betts\textsuperscript{6}, R.Bindel\textsuperscript{7}, W.Busza\textsuperscript{4}, Z.Chai\textsuperscript{2}, V.Chestlur\textsuperscript{9}, E.Garcia\textsuperscript{6}, T.Gbracek\textsuperscript{5}, K.Gulbrandsen\textsuperscript{4}, J.Hamblen\textsuperscript{8}, I.Harmar\textsuperscript{9}, C.Henderson\textsuperscript{4}, D.J.Hoffman\textsuperscript{6}, R.S.Hollis\textsuperscript{6}, R.Holyński\textsuperscript{3}, B.Holzman\textsuperscript{2}, A.Iordanova\textsuperscript{8}, J.L.Kane\textsuperscript{6}, P.Kulinich\textsuperscript{6}, C.M.Kuo\textsuperscript{5}, W.Li\textsuperscript{4}, W.T.Lin\textsuperscript{6}, C.Loizides\textsuperscript{8}, S.Manly\textsuperscript{8}, A.C.Mignerey\textsuperscript{7}, R.Nouicer\textsuperscript{2}, A.Olszewski\textsuperscript{3}, R.Pak\textsuperscript{2}, C.Reed\textsuperscript{4}, E.Richardson\textsuperscript{7}, C.Roland\textsuperscript{14}, G.Roland\textsuperscript{14}, J.Sagerer\textsuperscript{6}, I.Sedlak\textsuperscript{2}, I.Smith\textsuperscript{6}, M.A.Stankiewicz\textsuperscript{2}, P.Steinberg\textsuperscript{2}, G.S.F.Stephans\textsuperscript{4}, A.Sukhanov\textsuperscript{2}, A.Szostak\textsuperscript{2}, M.B.Tonjes\textsuperscript{7}, A.Trupe\textsuperscript{3}, G.J.van Nieuwenhuizen\textsuperscript{1}, S.S.Vaurynovich\textsuperscript{4}, R.Verdier\textsuperscript{4}, G.I.Veres\textsuperscript{4}, P.Walters\textsuperscript{8}, E.Wenger\textsuperscript{4}, D.Willhelm\textsuperscript{7}, F.L.H.Wolfs\textsuperscript{8}, B.Wosiak\textsuperscript{2}, K.Wozniak\textsuperscript{2}, S.Wyngaard\textsuperscript{2}, B.Wyslouch\textsuperscript{4}

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Elliptic flow, elliptic flow fluctuations and fluctuations in the initial geometry point to a description of nuclear collisions that is driven by the initial geometry, a quantity which appears to be imprinted from the instant of the collision. In these proceedings, recent results from the PHOBOS collaboration are discussed in the context of the importance of the collision geometry.

1 Introduction

Since the start of the RHIC program the measurement of particle azimuthal anisotropy, or flow, has been considered as one of the most important probes of nuclear collisions. Elliptic flow, in particular, is an important property of particle production as it is sensitive to the early stages
Figure 1: The left panel visualizes the two approaches to calculating eccentricity. The purple region (at center) in each collision illustrates the interacting nucleons. The orange and yellow nucleons (away from collision zone) are assumed not to directly influence the eccentricity. The solid (dashed) line represents the collision (participant) reaction plane. The lower part shows that the assumed reaction plane is rotated into the plane which maximizes the eccentricity, i.e., aligned along the semi-minor axis of the participant region. The right panel shows the difference of these two approaches for both Au+Au and Cu+Cu collisions. Cu+Cu collisions show a significant difference in the calculated eccentricity, whilst the discrepancy is less for Au+Au collisions.

of the collision and thus its study affords unique insights into the properties of the hot, dense matter that is produced in these collisions. At the root of this measurement lies a connection to the initial overlap geometry of the colliding nuclei, in particular the eccentricity of the initial overlap region of nucleons, which can be discussed as an averaged or event-by-event property of the system. The PHOBOS experiment has measured the elliptic flow for Au+Au and Cu+Cu collisions from $\sqrt{s_{NN}} = 19.6$ to 200 GeV, versus centrality and transverse momentum. For 200 GeV Au+Au collisions, a new analysis of the fluctuations in the magnitude of elliptic flow have revealed a startling agreement with a simple geometrical model of nuclear collisions.

2 Initial Collision Geometry

The collision geometry has always played an important role in heavy-ion collision analysis. The most simplistic definitions of centrality, derived from a Glauber model, and consequently the number of nucleons, $N_{\text{part}}$, expected to have participated in the collision is fundamental to this area of high-energy physics. As well as $N_{\text{part}}$, additional information can be gained from this model, including the spatial anisotropy of the collection of participating nucleons, or eccentricity ($\epsilon$). This anisotropy leads to the observed elliptic flow signal in data, discussed in the next sections. There are several methods for calculating $\epsilon$, two of which are illustrated in Fig. ??.

On the left, a schematic depiction of the “standard” (top, $\epsilon_{\text{std}}$) and “participant” (bottom, $\epsilon_{\text{part}}$) methods are shown. The former assumes that the collection of participating nucleons is oriented such that the semi-minor axis is aligned along the reaction plane - through the centers of the original colliding nuclei. As one can see, this is not always the case and may thus result in a reduced eccentricity. For the participant method, the semi-minor axis is allowed to rotate, such that the eccentricity is maximized. Eqn. ?? is a mathematical representation of the eccentricity for both methods.

$$
\epsilon_{\text{std}} = \frac{\sigma_y^2 - \sigma_x^2}{\sigma_y^2 + \sigma_x^2}
$$

$$
\epsilon_{\text{part}} = \frac{\sqrt{(\sigma_y^2 - \sigma_x^2)^2 + 4\sigma_x^2}}{\sigma_y^2 + \sigma_x^2}
$$

The difference in mean eccentricity between these two methods can be seen on the right panel of Fig. ???. For central Au+Au collisions little difference is observed between the two. For
more peripheral Au+Au or Cu+Cu collisions, large differences are seen, due primarily to the finite number of participating nucleon in such collisions. This difference in the magnitude of the eccentricity calculated using both methods from the model is observed in the elliptic flow data.

3 Elliptic Flow

Measurements of the elliptic flow, $v_2$, from PHOBOS are made over a broad range of pseudorapidity, centrality and energy. Generic features of particle production are found for both the Au+Au and Cu+Cu systems. At midrapidity, for similar centrality selections, the magnitude of $v_2$ increases from the lowest collision energy of $\sqrt{s_{NN}} = 19.6$ GeV up to 200 GeV. The magnitude of the $v_2$ diminishes as the pseudorapidity increases (for more forward particles) and is found to have a roughly triangular shape. The coupling of the collision energy and pseudorapidity dependences result in the $v_2$ signal exhibiting an extended longitudinal scaling behaviour, whereby the magnitude of $v_2$ is the same at the same pseudorapidity relative to beam rapidity (i.e. in the rest frame of one of the incoming nuclei).

The centrality dependence of $v_2$ shows the first clear dependence of the particle distributions following the underlying geometrical shape. For central Au+Au collisions with an almost full overlap (small impact parameter) both the $v_2$ and the eccentricity are found to be small, see Fig. ??a. As the impact parameter increases, collisions assume an almond shape, and $v_2$ and the eccentricity both increase.

For Au+Au collisions, it is found that $v_2$ scales reasonably with the standard eccentricity, $\epsilon_{std}$, whereas the Cu+Cu data strongly violate this approximate scaling, see Fig. ??a. Considering the alternate technique, the participant eccentricity, yields a unification of the two data samples, Fig. ??b.
4 Elliptic Flow Fluctuations

The collision species dependence of the integrated elliptic flow signal is found to be strongly
dependent on the collision geometry, and to its precise definition. Specifically, the fluctuations
in the nucleon positions on an event-by-event basis appears to drive the final $v_2$ signal. If
such fluctuations influence the averaged signal, then this should be a measurable quantity in
itself. One of the latest results from the PHOBOS collaboration concentrates on measuring
these elliptic flow fluctuations. The method utilizes the whole pseudorapidity coverage of the
PHOBOS detector to measure the $v_2$ signal on an event-by-event basis, assuming the shape is
either a triangle or a trapezoid. Details of the analysis method can be found in Ref. 7.

The elliptic flow fluctuations, expressed as $\sigma_{v_2}/v_2$, are shown in Fig. 7. The fluctuations
are found to be significant for all centrality classes studied, with a peak close to 50% relative
fluctuations. Fluctuations in the eccentricity from the Glauber model calculations are also found
to be significant, with the magnitude in remarkable agreement with the $v_2$ fluctuations. Such
an agreement hints that the detailed initial geometrical configuration is imprinted on the final
distribution of particles.

5 Summary

The initial geometry in nuclear collisions plays an important role in particle production at RHIC.
The detailed eccentricity, calculated from the positions of the interacting nucleons in a Glauber
model, has been shown to unify elliptic flow data from Au+Au and Cu+Cu collisions. The
magnitude of elliptic flow fluctuations are measured and are found to be large for all centralities.
The level of these fluctuations is strikingly similar to those from the eccentricity calculations,
driving that the initial geometry is imprinted on the final particle distributions.

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Systematics of Soft Particle Production at RHIC: Lessons from PHOBOS

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The PHOBOS experiment has measured the properties of particle production in heavy ion collisions between \( \sqrt{s_{NN}} \) of 20 and 200 GeV. The dependencies of charged particle yield on energy, system size, and both longitudinal and transverse momentum have been determined over the full kinematic range. Identified charged particles emitted near mid-rapidity have been studied over about 2 orders of magnitude in transverse momentum. This broad data set was found to be characterized by a small number of simple scalings which factorize to a surprising degree. This study has recently been extended by the addition of new data for Cu+Cu as well as new analyses of Au+Au data, including more peripheral collisions. In addition, the exploration of global properties has been expanded with the use of new techniques, including two-particle correlations, more sensitive searches for rare events, and more detailed studies of particles emitted at very forward rapidity. The characteristics of particle production which are revealed by this extensive data set will be described along with the implications for future data from the LHC.

1 Introduction

The PHOBOS experiment took data at RHIC starting with the first beam in June of 2000 and continuing through Run 5 in the spring of 2005. Data were taken for a broad range of systems, namely p+p at two energies, d+Au at one energy, Cu+Cu at three energies, and Au+Au at five energies. Results from PHOBOS and the three other RHIC experiments have shown that heavy ion collisions at the highest RHIC energies result in the formation of a new state of matter, characterized by a high energy density and dominated by partonic degrees of freedom.\textsuperscript{1}

One of the primary goals of the PHOBOS experimental program was the characterization
of the properties of particle production over a very broad range in energy and system size, as well as over several orders of magnitude in transverse momentum and all or a very large fraction of the pseudorapidity distribution. While not necessarily evidence for, or a direct probe of, the exotic partonic state, these observables set constraints on models of the formation and subsequent hadronization of the novel medium. Final state particle distributions can also set limits on basic properties of the system such as energy density and entropy. In addition to contributing significantly to our understanding of the systems formed at RHIC, this extensive data set has revealed a number of surprising results.

2 Energy Dependence of Particle Production

The first physics result from RHIC was the PHOBOS publication of the charged particle pseudorapidity density, $dN/d\eta$, near mid-rapidity at nucleon nucleon center of mass energies ($\sqrt{s_{NN}}$) of 56 and 130 GeV. This early result had three immediate and profound impacts on the field of relativistic heavy ions. First, the numerical value invalidated the majority of the theoretical predictions in existence at the time. Second, the data lent support to concepts of parton saturation, which if validated could describe the dominant physics process controlling the low-x region at high energies, even in p+p collisions. Finally, the fact that these values were significantly lower than many of the theoretical predictions, combined with the first PHOBOS $dN/d\eta$ data at $\sqrt{s_{NN}}=200$ GeV which was also on the low side of the revised theoretical predictions, suggested that tracking and other measurements in heavy ion collisions at the LHC might not be as formidable as originally thought. This realization helped to spawn a significant expansion in the planned LHC heavy ion program.

Later analysis revealed an intriguing similarity between the particle multiplicities per pair of participating nucleons in nucleus-nucleus collisions when compared to proton-(anti)proton interactions at twice the center of mass energy (i.e. $(p+p)/\sqrt{s}=2\times(A+A)/\sqrt{s_{NN}}$). Further, these multiplicities were similar to those seen in $e^+e^-$ at the same energy. The comparison of p+p and $e^+e^-$ was known previously and assumed to be due to the fact that only about half of the center of mass energy in p+p was available for particle production. The new comparison with Au+Au implies that nucleus-nucleus collisions can convert a much larger fraction of the available energy into particles. Data from the LHC (p+p at 14 TeV and Pb+Pb at 5.5 TeV) will reveal whether or not this correspondence extends to much higher energies.

3 Pseudorapidity Dependence of Particle Production

The uniquely broad pseudorapidity coverage of the PHOBOS multiplicity detector allowed measurement of all or almost all of the $dN/d\eta$ distribution, even at the highest RHIC energies. In addition to producing total multiplicity data with relatively small systematic errors, these results made possible a detailed comparison of the shape of the distribution at different center of mass energies. When the $dN/d\eta$ distributions for Au+Au ranging from $\sqrt{s_{NN}}=19.6$ to 200 GeV were plotted as a function of $|\eta−\eta_{beam}|$, thereby effectively viewing them in the rest frame of one of the colliding nuclei, it was found that data from all energies followed a common curve (see top left panel of Fig. 1). Furthermore, preliminary data for Cu+Cu over roughly the same range in energy reveal that they follow exactly the same curve. Thus, the “limiting fragmentation” or “extended longitudinal” scaling seen previously in small systems was found to apply also in heavy ion collisions. Assuming that this observation and the energy dependence described in the previous section extend up to LHC energies, an empirical prediction of the full $dN/d\eta$ shape can be made.

A related, but much more surprising, result was found when a similar analysis was applied to the pseudorapidity dependence of elliptic flow. When the data for $v_2$ from Au+Au were plotted in the effective rest frame of one of the colliding nuclei, a pattern of “extended longi-
Figure 1: (Top left) Charged particle $dN/d\eta$ for various systems and energies effectively viewed in the rest frame of one of the colliding nuclei. (Bottom left) A similar plot for elliptic flow. (Right) $dN/d\eta$ per participant pair versus centrality for Au+Au at four energies fit using a product of separate functions of energy and centrality.

The transverse scaling was again revealed (see bottom left panel of Fig. 1). This adds an additional intriguing element to the experimental results for the pseudorapidity dependence of elliptic flow, data which have presented a considerable challenge to existing theories (see for example). Again, an extrapolation to LHC energies can be done but the interpretation of the result is unclear. The magnitude of $v_2$ at midrapidity in Au+Au at $\sqrt{s_{NN}}=200$ GeV is claimed to saturate the hydrodynamical limit but, if the observed trend continues, the value at the LHC will be significantly larger.

4 Centrality Dependence of Particle Production

By analyzing heavy ion collisions at varying impact parameter, the effect of system size on particle production can be explored. This variation does not represent simply “more of the same” since central collisions with small impact parameters have a larger average number of collisions per participant and therefore might differ more significantly from elementary $p+p$ interactions. A common claim is that “harder” processes should scale with the number of collisions while “softer” processes should scale with the number of participants. If true, this belief, combined with the expectation that the ratio of “hard” over “soft” processes should increase with collision energy, implies that the centrality dependence must be energy dependent. In stark contrast to this prediction, the centrality dependence is found to be identical at all energies studied. In fact, it can be shown quantitatively that the data factorize by fitting them with the product of separate functions of energy and the number of participating nucleons (see right panel of Fig. 1). Far from being a property solely of bulk particle production at lower transverse momentum, this factorization was found to extend up to $p_T$ of almost 4 GeV/c in Au+Au data. Again, Cu+Cu results are observed to follow the same trend. As with many of the PHOBOS observations, it will be very interesting to follow this trend to the LHC where “harder” processes are expected to
make a much more dominant contribution to particle production.

5 Continuing PHOBOS Analysis

Although no further data are being taken by the PHOBOS collaboration, analysis work continues. One goal is to fully incorporate all results, including those for smaller systems such as p+p and nucleus-nucleus collisions over an extended range of centrality. Simultaneously, the analysis is expanding beyond event-integrated single-particle distributions to the consideration of more complicated observables such as fluctuations, correlations, and rare event topologies. Results for elementary systems, to be used as a baseline comparison for nucleus-nucleus data, have already been published.15

6 Summary

Analysis of the characteristics of particle production in nucleus-nucleus collisions at RHIC energies have revealed a number of unexpectedly simple dependencies. Observables considered range from the most basic such as total multiplicity to the fairly complex such as elliptic flow. In many cases, the dependencies on collision energy, centrality, pseudorapidity, and transverse momentum factorize to a surprising degree. To paraphrase a comment originally made about star formation in galaxies16, “Particle production in heavy ion collisions follows a quite simple pattern and simple patterns often mean that there are only a few basic physical mechanisms at work. . . . We can now find out what these mechanisms are by measuring how particle production behaves with energy, centrality, \( \eta \), and \( p_T \) and compare that behavior to models”. Extrapolation of these trends to LHC energies suggest that interesting discoveries may well be made using only these simply global observables.

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HBT with an emphasis on exotic particle femtoscopy

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STAR latest analyses of two-particle correlations are presented, including measurements with heavy, exotic, strange and multi-strange hadrons (p,K,Λ,Ξ). These measurements yield insight into the dynamics of heavy-ion collisions. Species-dependent emission sizes and emission asymmetries are expected to arise as a consequence of collective expansion. Furthermore the correlation analyses in heavy-ion collisions allow to study, otherwise hard to measure, strong interaction potential between different hadrons.

Keywords: heavy-ion, RHIC, STAR, femtoscopy, HBT

1 Introduction

Heavy-ion experiments have been pursuing the study of nuclear matter under the extreme conditions created in the collisions of large nuclei at high energies with a motivation to observe phase-transition into de-confined partonic matter as it was predicted by QCD calculations. Recent results from RHIC\(^1\) indeed indicate that strongly interacting matter governed by partonic degrees of freedom was created. Measurements reveal that system undergoes rapid collective expansion with possible early thermalization of the constituents. The non-trivial space-time evolution of the system before and also after return into the state of confined nuclear matter is one of the main features of the heavy-ion physics.

Two-particle correlation femtoscopy\(^2\) (often called HBT) is one of most direct ways of studying the space-time structure of the particle-emitting source and can shed a light onto the dynamics of the collision. Particles emitted with small relative momenta are correlated due to their final state interaction (FSI) and/or their interference due to quantum statistics (QS). The
two-particle correlation function is defined as:

$$C^{ab}_{R}(\vec{q}) = \frac{d^{3}N_{a}^{ab}/(dp_{a}^{3}dp_{b}^{3})}{(d^{3}N_{a}^{a}/dp_{a}^{3})(d^{3}N_{b}^{b}/dp_{b}^{3})} = \int d^{3}q^{a} \cdot S_{R}^{ab}(\vec{q}) \cdot |f(\vec{q}, \vec{r})|^{2},$$

where $\vec{K} = \vec{p}_{a}^{b} - \vec{p}_{b}^{a}$, $\vec{q} = \vec{p}_{a} - \vec{p}_{b}$ and $f(\vec{q}, \vec{r})$ is the pair relative wavefunction, including the FSI and QS describing the propagation of the pair from a relative separation of $\vec{f}$ to the detector with relative momentum $\vec{q}$. The source function $S_{R}^{ab}(\vec{q})$ is the probability of emitting a pair of particles with average momentum $\vec{K}$ at a distance $\vec{r}$ apart. Hence the correlation function (1) encodes the information about space-time configuration of the emitting source. It must be stressed, however, that $S_{R}^{ab}(\vec{q})$ is not sensitive to the size of the entire source, but to the so-called “homogeneity region” - part of the phase space occupied by outgoing particles whose velocities have a specific magnitude and direction.

2 Expansion dynamics and HBT observables

When the medium undergoes expansion it means that there is strong correlation between momenta of outgoing particles and their emission points. As a consequence, such models as hydro
and microscopic models, which include significant collective expansion, predict decrease of the measured radii (i.e. the size of the homogeneity region) with increasing transverse mass $m_t$ of the particles. In Figure 1 is shown a collection of results on measured pion source sizes by different experiments for their largest colliding system and for the most central collisions. The $m_t = \sqrt{K_t^2 + m_0^2}$ dependence of HBT radii is generally attributed to the collective expansion. Further measurements show that not only pions but also other particle species including strange and multi-strange particles significantly participate in the collective expansion. Thus the femtoscopic signals arising from it, such as the fall of pion HBT radii with $m_t$, should also be observed. Recent high statistics data collected at RHIC has allowed to carry out femtoscopic analyses with precision that was unreachable before and for particle pair types which were measured for the first time. In Figure 2 is presented $m_t$ dependence of radii for different measured pairs of particles with close mass. It is of particular importance to note that these measurements include systems with different final state interaction and different systematical uncertainties. Using different particle species one can “turn on/off” Coulomb FSI and/or quantum statistics effects. In case of $K^0_S - K^0_S$ and $p(\bar{p}) - \Lambda(\bar{\Lambda})$ the Coulomb interaction is absent and in the latter case there is also no QS. The $p - \Lambda$ and $p - \bar{\Lambda}$ analyses are an example of a way hadron interactions can be studied using femtoscopy, as presented by the STAR experiment in. While in $p - \Lambda$ and $p - \bar{\Lambda}$ strong interaction is known, allowing to perform standard femtoscopic measurement of the source size, $p - \Lambda$ and $p - \Lambda$ was measured for the first time and the interaction is unknown. However, assuming the same functional form of the interaction as in $p - \Lambda$, $p - \bar{\Lambda}$ and treating the potential parameters (scattering lengths) as free parameters it was possible to extract spin-averaged scattering lengths as presented in Figure 3. The message from common scaling in Figure 2 is strengthened when considering different systematics involved in each measurement. From this point $p(\bar{p}) - p(\bar{p})$ analyses are especially of interest since these measurements are strongly affected by residual correlations originating from decays of $\Sigma$ and $\Lambda$. This issue has been extensively treated in.

Femtoscopic measurements with non-identical particles are not only sensitive to the size of the system, but also to the relative difference in the space-time position of the emission of the two particle species. Models that include collective expansion predict a relation between the average emission position and the mass of the particle such as that particles with higher $m_t$ are emitted more on the outside of the expanding fireball. This effect hence increases with a mass difference within the measured particle pair. Figure 4 shows that in heavy-ion collisions the
average emission points of particles with different mass, such as pions, kaons and protons, are significantly shifted with respect to each other. These measurements were recently extended by STAR experiment to include multi-strange baryons using $\pi - \Xi$ correlations. This exotic system is of particular interest as it includes an order of magnitude difference in mass plus $\Delta B = 1$ and $\Delta S = 2$ gap in baryon and strangeness quantum numbers respectively. The case of multi-strange baryon flow is of high importance since the collective behavior of these particles, suggested by large values of observed elliptic flow, is believed to be coming predominantly from the early partonic stage. Recent results on $\pi - \Xi$ correlation function for different combination of charges are presented in Figure 5, where $2k = \bar{q}$ in the pair cms. Spherical decomposition method, which has recently become a promising tool for 3-dimensional analyses of the correlation function, is used. As described in, the non-zero value of AT-coefficient signalizes space-time shift in the average emission between pions and $\Xi$s. Fit of the Coulomb interaction to the data yields value of the space-time shift $(5.6 \pm 1.)$ fm, including statistical errors only. This significantly large value is also in qualitative agreement with the collision evolution during which multi-strange baryons take part in the collective expansion of the matter.

3 Conclusions

STAR measurements on two-particle femtoscopy were presented, emphasizing results with heavy, strange and multi-strange hadrons. A species-independence of the m scaling of HBT radii together with extracted emission asymmetries among the particles were shown. These results provide an independent confirmation of a transversely expanding particle source in heavy-ion collisions. Results on $p - \Lambda$, $p - \Lambda$ show that non-identical correlations can be used to study otherwise hardly accessible hadron interactions.

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Intra-jet correlations of high-$p_t$ hadrons from STAR

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Systematic measurements of pseudorapidity ($\Delta\eta$) and azimuthal ($\Delta\phi$) correlations between high-$p_t$ charged hadrons in $\sqrt{s_{NN}}=200$ GeV Au+Au collisions are presented. An enhancement of correlated yield at large $\Delta\eta$ on the near-side is observed (the ridge). This effect persists up to trigger $p_{t,\text{trig}}^{\text{trig}} \sim 9$ GeV/c, indicating that it is associated with jet production. More detailed analysis suggests that the near-side $\Delta\eta \times \Delta\phi$ correlation consists of two distinct components: a jet contribution and the ridge contribution with properties closely related to the medium produced in central Au+Au collisions.

Dihadron azimuthal correlation studies in nuclear collisions have shown that hard partons interact strongly with the matter that is generated and provide a sensitive probe of the medium. Enhanced near-side ($\Delta\phi \sim 0$) correlated yield at large $\Delta\eta$ (the ridge) has been observed in measurements with trigger particles at intermediate $p_t$ ($4 < p_{t,\text{trig}}^{\text{trig}} < 6$ GeV) and for dihadron pairs having $p_t < 2$ GeV but no trigger requirement. However, inclusive hadron production at $p_t \leq 6$ GeV/c exhibits large differences between nuclear collisions and more elementary collisions. It is therefore unclear from these existing measurements whether the ridge is associated with hard partonic scattering and jet production. In this proceeding we extend the near-side correlation measurement to $p_{t,\text{trig}}^{\text{trig}} \sim 9$ GeV/c, well into the kinematic region where inclusive hadron production is similar in nuclear and elementary collisions and where jet fragmentation is thought to dominate. We observe the persistence of the ridge effect to the highest measured trigger $p_t$, suggesting that it is indeed associated with jet production. We further characterize the ridge, to gain insights into its origin.

To illustrate the analysis method, Fig. 1 shows the $\Delta\eta \times \Delta\phi$ distribution of hadrons with $p_{t,\text{assoc}}^{\text{assoc}} > 2$ GeV, associated with trigger hadrons $3 < p_{t,\text{trig}}^{\text{trig}} < 4$ GeV in central Au+Au collisions. The yields are corrected for single-particle tracking efficiency and the finite pair acceptance in $\Delta\eta$ and $\Delta\phi$. The near-side yield shows a clear peak around ($\Delta\eta, \Delta\phi$) = (0,0), as expected from jet fragmentation. In addition a prominent enhancement of correlated yield in $\Delta\eta$ on the near-side is clearly visible above the flow modulated background (the ridge). To better understand the ridge phenomenon we decompose the near-side into a jet-like peak and a $\Delta\eta$ independent ridge component. This ansatz assumes distinct underlying physical processes in the different $\Delta\eta$ regions. We examine this assumption below.

To extract the ridge yield from dihadron measurements we project the two dimensional ($\Delta\eta \times \Delta\phi$) correlation function (Fig. 1) onto $\Delta\phi$ and $\Delta\eta$ in different $\Delta\eta \times \Delta\phi$ regions. Three methods were used to characterize the small $\Delta\eta$ jet-like ($J$) and the large $\Delta\eta$ ridge-like ($R$) contributions to the near-side jet yield in $\Delta\eta$ and $\Delta\phi$:

\[a\] The yields are extracted from bin-counting in the interval $|\Delta\phi| = |\Delta\eta| < 1$
Figure 1: Raw $\Delta \eta \times \Delta \phi$ dihadron correlation function in central Au+Au collisions for $3 < p_t^{\text{trig}} < 4$ GeV and $p_t^{\text{assoc}} > 2$ GeV.

- $\Delta \phi(J + R)$: Projecting onto $\Delta \phi$ with the full experimental $\Delta \eta$ acceptance ($|\Delta \eta| < 1.7$ was used in this analysis) and subtracting the elliptic flow ($v_2$) modulated background.

- $\Delta \phi(J)$: Subtracting the $\Delta \phi$ projection for $0.7 < |\Delta \eta| < 1.4$ from the $\Delta \phi$ projection $|\Delta \eta| \leq 0.7$ (near-side).

- $\Delta \eta(J)$: Projecting onto $\Delta \eta$ in a $\Delta \phi$ window $|\Delta \phi| < 0.7$ (near-side). A constant fit to the measurements was used to subtract the background.

In Fig. 2 the near-side yield is shown as a function of the number of participants $N_{\text{part}}$ for all three methods. The agreement of the measured jet-like yield between the $\Delta \eta(J)$ and $\Delta \phi(J)$ method for all centrality bins, within the sensitivity of this analysis, supports the assumption that the ridge-like correlation is uniform in the $\Delta \eta$ acceptance. Further detailed studies of the ridge shape in the high statistics Au+Au central data set, especially at high $p_t^{\text{trig}}$, will be pursued. Note that the jet-like correlated yield is independent of centrality and agrees with the p+p reference measurements. In contrast the $\Delta \phi(J + R)$ yield shows a significant increase with centrality due to the inclusion of the correlated yield at large $\Delta \eta$ (ridge).

For the purpose of this analysis one can define the (absolute) ridge yield $= \text{yield}(\Delta \phi(J + R)) - \text{yield}(\Delta \eta(J))$. The main systematic error is the uncertainty in the elliptic flow measurement for the $\Delta \phi(J + R)$ method. The $v_2$ value used in this analysis is the mean of the reaction plane ($v_2\{RP\}$) and four-particle cumulant method ($v_2\{4\}$) in Au+Au collisions. The systematic uncertainties were estimated using $v_2\{RP\}$ as maximum and $v_2\{4\}$ as minimum $v_2$ values (represented as lines in all figures).

Fig. 3 shows that a significant (absolute) ridge yield persists up to the highest $p_t^{\text{trig}}$, with yield increasing with centrality. The finite ridge yield at $p_t^{\text{trig}}$ up to 9 GeV, where parton fragmentation is expected to be the dominant hadron production mechanism, indicates that the ridge is associated with jet production.

To characterize in more detail the properties of particles associated to the ridge-like or jet-like near-side correlation we use the $p_t^{\text{assoc}}$ spectrum in different $p_t^{\text{trig}}$ windows, as shown in Fig. 4. An exponential function $\frac{dN}{dp_t^{\text{assoc}}} \propto p_t e^{-p_t/T}$ is fitted to the data (lines in Fig. 4) to extract the inverse slope parameter $T$. A clear difference between the slopes of the jet-like yield, using the $\Delta \eta(J)$ method, and the ridge-like yield is seen: while the $p_t$ dependence of the ridge yield

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*Not corrected for the finite $\Delta \eta$ pair-acceptance.

*The ridge yield depends on the $\Delta \eta$ integration window used in the $\Delta \phi(J + R)$ method.
is similar to the inclusive particle production\textsuperscript{10}, the jet-like associated yield has a significantly harder $p_T$-spectrum, increasing with $p_T^{\text{frag}}$, as expected from jet fragmentation. The slope of the ridge-like yield is largely independent of $p_T^{\text{frag}}$ and only slightly harder than the inclusive spectrum with a slope difference $\Delta T \approx 40-50$ MeV.

Fig. 5 a) shows the $p_T^{\text{frag}}$ dependence of the near-side $z_T$ di-hadron fragmentation function ($z_T = p_T^{\text{assoc}}/p_T^{\text{frag}}$) in central Au+Au collisions (for details see\textsuperscript{11}). Subtracting the ridge-like contributions, using the $\Delta \phi(J)$ method, one observes a near-side fragmentation that is approximately independent of $p_T^{\text{frag}}$ (Fig. 5 b)). The $z_T$ distributions in central Au+Au collisions after subtracting the ridge contribution are comparable to the d+Au reference measurements (Fig. 5 c)) in contrast to the non-ridge subtracted distributions\textsuperscript{11}. To further quantify this observation, studies to estimate the effect of background fluctuations in the $\Delta \phi(J)$ method at high $\Delta \eta$ will be pursued.

These observations support the ansatz that the near-side $\Delta \eta \times \Delta \phi$ correlation consists of two distinct components: a jet contribution, consistent with the $p+p$ and d+Au dihadron reference measurements\textsuperscript{4,5}, and the ridge contribution with properties similar to the medium. This could arise from partonic energy loss followed by fragmentation in vacuum, with the lost energy appearing dominantly in the ridge.

Several models are qualitatively able to describe the presented phenomena: coupling of induced radiation to longitudinal flow\textsuperscript{12}, turbulent color fields\textsuperscript{13}, anisotropic plasma\textsuperscript{14}, a combination of jet-quenching and strong radial flow\textsuperscript{15} and recombination of locally thermal enhanced partons due to partonic energy loss in the recombination framework\textsuperscript{16}. A comparison of quantitative theoretical calculations to the measurements are needed in order to understand the origin of the ridge.

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Figure 5: (color online) Near-side $z_T$ dihadron fragmentation function for different $p_T^{rig}$ in central Au+Au collisions before a) and after ridge subtraction b) as well as the ratio to d+Au reference measurements c).

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MEDIUM-MODIFIED FRAGMENTATION FUNCTIONS

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We extend the Sudakov formalism from vacuum to opaque media by supplementing the splitting functions with an additional term given by the medium-induced gluon radiation spectrum. We then solve the DGLAP evolution equations to obtain the medium-modified fragmentation functions. In this way, both the additional energy loss and the modification of the QCD evolution by the medium are included in a consistent manner. As a phenomenological application, we compute the suppression of the high-$p_t$ yields in heavy-ion collisions by convoluting the obtained fragmentation functions with the perturbative spectrum.

Keywords: Jet Quenching, Heavy Ion Collisions, Jets

1 Modification of the vacuum splitting functions

A hard parton produced in a heavy ion collision will travel through the opaque medium losing virtuality until it eventually hadronizes, which takes place in vacuum for large enough transverse momentum. The medium accelerates the QCD evolution of the parton by inducing successive soft gluon radiations. Under the BDMPS approximation, the spectrum of gluons emitted by the parent parton can be written as a function of two parameters which completely characterize the medium: the medium length $L$ and the transport coefficient $\hat{q}$. The properties of the medium are encoded in the transport coefficient $\hat{q}$, which can be related to the average squared transverse momentum transferred to the parton per mean free path.

To compute medium modified fragmentation functions (MMFF), we extend the fact that the formalism provides the vacuum splitting functions in the collinear limit, $dP^{VAC}/dzdk_T = \alpha_s P(z)VAC/2\pi k_T^2$, to the medium, $dP^{MED}/dzdk_T = \alpha_s P(z)^{MED}/2\pi k_T^2$. The total splitting function is the sum of both (see also),

\[ P_{5a} = P_{5a}^{VAC}(z) + P_{5a}^{MED}(z, t, \hat{q}, L). \]

See that eq. 1 is only valid at $z \to 1$ due to the use of the medium induced gluon spectra, which relies on high energy approximations. So, at this level, the only difference between our total $g \to gg$ and $q \to qg$ splittings will be still a Casimir color factor. To extend it to all $z$, we take the full LO splitting functions $P^{VAC}$. Small $z$ corrections are also introduced in $P^{MED}$, whose effect is found to be small (see fig. 1 below).
2 Sudakov Factors

The Sudakov form factor is defined in the standard way

\[ \Delta_a(t, t_0) = \exp - \sum_{a \to ac'} \int_{t_0}^{t} \frac{dt'}{t'} \int_{z_{min}(t')}^{1} \frac{dz}{2\pi} \frac{\alpha_s(z(1-z)t')}{\alpha_s(z(1-z)t')} P_{ac}(z), \]

which can be interpreted as the probability for a parton a not to radiate resolvable partons when evolving between the scales \( t_0 \) and \( t \). In fig. 1 we show the effect of supplementing the Sudakov form factor with a medium term in the splitting function as given by eq. (1). This is done for two energies \( E = 10 \text{ GeV} \), of interest for RHIC, and \( E = 100 \text{ GeV} \), of interest for the LHC. Increasing the value of the transport coefficient leads to an enhancement on the radiation and, hence, to more suppressed Sudakov factors. The Sudakov for gluons is approximately enhanced by the color factor \( 9/4 \) with respect to that for quarks.

3 Medium-Modified Fragmentation Functions

DGLAP evolution can be written in terms of the Sudakov factors:

\[ \frac{D_k^h(x, t)}{\Delta_a(t_0, t)} = D_k^h(x, t_0) + \int_{t_0}^{t} \frac{dt'}{t'} \frac{1}{\Delta_a(t_0, t')} \int_{z_{min}(t')}^{1} \frac{dz}{2\pi} \frac{\alpha_s(z(1-z)t')}{\alpha_s(z(1-z)t')} \sum_b P_{ba} D_b^h \left( \frac{x}{z}, t' \right), \]

so that the probability that a parton a fragments into a hadron at the scale \( t \) with an energy fraction \( x \) is equal to the probability that the parton has not radiated since it was produced at \( t_0 \) (first term) plus the probability of having radiated at any intermediate scale \( t' \) (second term). We consider only tree flavors (u, d, s), we use as initial conditions the KKP fragmentation functions\(^a\) at \( t_0 = 2 \text{ GeV}^2 \) and the scale in \( \alpha_s \) is taken to be the transverse momentum of the emission \( Q^2 = z(1-z)t \). The virtuality is evolved between \( t_0 < t < 4E^2 \) and the infrared cut is chosen as \( t_0/2t < z(t) < 1 - t_0/2t \). For the vacuum, we have checked that this method reproduces the KKP results better than 40% in the \( (z, Q^2) \) region of interest.

In fig. 2 we show our results for the MMFF with the modified DGLAP evolution (3), for different parton energies and medium parameters, and for different parton types. A clear suppression at large \( z \) and enhancement at small \( z \) can be seen. These two characteristics grow
Figure 2: Fragmentation functions for gluons and quarks of different energies and for different medium characteristics (see the labels on the plots).

with increasing evolution, medium length and transport coefficient and, as expected, they are more pronounced for gluons than for quarks. The fact that these effects are larger for smaller parent parton energies can be understood qualitatively from the fact that the energy loss becomes more and more energy-independent with increasing energy.

Previous calculations of the MMFF were based on the quenching weights. This method had some limitations concerning energy and momentum conservation. The MMFF were calculated shifting the vacuum ones and there was no evolution in virtuality, nor was there a unified description of medium and vacuum. Within our new approach this limitations are overtaken: we modify the QCD evolution and vacuum and medium are evolved together.

4 Nuclear Modification factor

We convolute our MMFF with the nuclear parton densities (pdf) and the hard scattering elements to obtain medium modified particle spectra following the factorization formula at LO:

$$\sigma^{A_{c,t} h} = f_A(x_1, Q^2) f_B(x_2, Q^2) \otimes \sigma(x_1, x_2, Q^2) \otimes D_{h \rightarrow A}(z, Q^2).$$  \hspace{1cm} (4)

Our proton-proton reference is shown in fig. 3(left). Then we show our results for the nuclear modification factor(right) (defined as $R_{AA} = \frac{d\sigma}{d^3p_T}(pdf + EKS + MMFF) / \frac{d\sigma}{d^3p_T}(pdf + VACFF)$ ) in Au-Au collisions at a fixed length of $L = 6$ fm. All the three scales (factorization, renormalization, fragmentation) are set equal to the fragmenting parton momentum. The value
of the transport coefficient which better describes the data is $\hat{q} \sim 1 \text{ GeV}^2/\text{fm}$. This value is in agreement with the findings in the first paper in 6, where also a fixed medium length was used. Once a realistic geometry is taken into account, this value is known to increase. We do not attempt here to make a fit to experimental results but just to check that this new procedure, which explicitly conserves energy momentum at each splitting, results on values of $\hat{q}$ compatible with previous calculations.

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HIGHLIGHTS FROM THE NA60 EXPERIMENT

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The NA60 experiment is a fixed-target experiment at the CERN SPS. It has measured the dimuon yield in In-In collisions with an In beam of 158 AGeV/c and in p-A collisions with a proton beam of 400 and 158 AGeV/c. The results allow to address three important physics topics, namely the study of the ρ spectral function in nuclear collisions, the clarification of the origin of the dimuon excess measured by NA50 in the intermediate mass range, and the J/ψ suppression pattern in a collision system different from Pb-Pb. An overview of these results will be given in this paper.

The NA60 experiment. The measurement of dimuon production is a key tool to gain insight into ultra-relativistic nuclear collisions. However, the importance of the dimuon data depends strongly on the resolution of the experimental apparatus. While the NA50 experiment measured the J/ψ suppression pattern as a function of centrality in Pb-Pb collisions, its mass and vertexing resolutions were not sufficient to address two other important physics topics, namely the shape of the in-medium spectral function of the ρ meson and the origin of the dimuon excess observed by NA50 in the intermediate mass region (IMR, 1.2-2.7 GeV/c²).

The NA60 experiment inherited from NA50 the muon spectrometer (MS) for muon triggering and tracking, and the Zero Degree Calorimeter (ZDC), which measures the energy carried forward (at 0°) by the spectator nucleons, to evaluate the centrality of A-A events. Two components were added: before the target, a hemiscope (HS) made of two pairs of cryogenic silicon strip detectors determines the beam impact point on the target transverse plane with a 20 μm resolution; after the target, a vertex tracker (VT), made of 16 planes of radiation-hard silicon pixel detectors embedded in a 2.5 T dipolar magnetic field, measures tracks and their momenta before they suffer multiple scattering in the absorber (for details on the NA60 apparatus see 1).

Tracks in the VT are matched (both in momenta and in coordinate space) to tracks measured in the muon spectrometer. This procedure greatly improves the mass resolution, particularly in the low-mass range (20 MeV/c² at the ω mass, to be confronted with NA50 80 MeV/c²). This allows to resolve and evaluate the different contributions to the dimuon spectrum below 1.2 GeV/c. Moreover, it is possible to determine the offset between muon tracks and the primary interaction vertex with a precision of ~ 40 μm for 20 GeV/c muons. This is sufficient to discriminate between muons coming directly from the fireball and muons originated by secondaries decays in the intermediate mass range.

The results reported in this paper were obtained from the analysis of data taken in 2003 for In-In collisions at 158 AGeV/c, and in 2004 for p-A collisions at 158 GeV/c. The choice of a colliding system different from Pb-Pb makes it possible to search for the scaling variable which
drives the onset of the $J/\psi$ anomalous suppression. Moreover, p–A data at 158 GeV allow to compare p–A to A–A data without systematic errors deriving from the energy rescaling.

The low mass region. The net opposite-sign dimuon spectrum below 1.4 GeV/c\(^2\) has been obtained from the raw mass spectrum by subtraction of combinatorial background and of signal fake matches between tracks in the MS and in the VT. Four centrality classes were defined via the charged particle multiplicity measured by the VT. Most peripheral data ($4 < dN_{ch}/d\eta < 30$) are well reproduced by the cocktail of expected electromagnetic decays of the neutral mesons, while for more central collisions a strong excess appears, whose shape is not known a priori. Thanks to the high data quality, it is possible to isolate this excess by subtracting from the data a hadron cocktail without the $\rho$ ($s < 2$), whose spectral shape is expected to be modified in the fireball. The resulting excess is shown in Fig.1 (left) for semicentral collisions ($110 < dN_{ch}/d\eta < 170$): it is characterized by a peaked structure centered on the nominal $\rho$ mass, which broadens and increases with centrality and resides on a wide continuum. However, a quantitative analysis of the excess (for details see \cite{3}) demonstrates that the ratio between the continuum-subtracted peak and the $\rho$ obtained from the cocktail fit decreases by almost a factor 2 from most peripheral to most central bin. This means that the excess can not be simply interpreted as the cocktail $\rho$ on top of a continuum.

As shown in fig.1 (left), the moving mass model related to Brown/Rho (BR) scaling is ruled out. The qualitative features of the excess mass spectra are consistent with the interpretation as direct thermal radiation from the fireball, dominated by $\pi\pi$ annihilation. Models based on the in-medium $\rho$ broadening scenario and which take into account the role of baryons in the broadening\cite{4} are able to reproduce quantitatively the data below 0.9 GeV/c. For the mass region above 0.9 GeV, data seem to be described equally well by introducing $4\pi$ hadronic processes sensitive to vector-axialvector mixing (and therefore to chiral symmetry restoration) or in terms of partonic processes dominated by $q\bar{q}$ annihilation. This feature could be a manifestation of parton-hadron duality.

The study of $p_T$ spectra presents further interesting features. The trend at small $p_T$ is opposite to the flattening expected from radial flow of hadrons produced at kinetic freeze-out, while it flattens as expected in the $\phi$ mass bin. Moreover, the $T_{eff}$, obtained from fits of the
$m_T$ spectra (see fig. 1 right) and plotted as a function of mass, shows a maximum in the $\rho$-like region. This may be attributed to the $\rho$ produced at freeze-out, which experiences the largest blue-shift and thus the highest effective temperature. Besides that, it is worth noting that the continuum above 0.9 GeV/c is cooler than the continuum below 0.6 GeV/c; this seems to indicate that the two regions are fed by qualitatively different sources. It is hoped that a finer theoretical understanding of $p_T$ spectra could serve as a handle to disentangle partonic from hadronic sources (breaking parton-hadron duality). For details see 6.

Intermediate mass region. To understand the origin of the excess measured by NA50 in the IMR, NA60 has measured the offset of the muon tracks with respect to the main interaction vertex. This distribution has been weighted by the inverted error matrix from the vertex fit and the muon extrapolation (the combinatorial background was subtracted by event mixing). The resulting distribution is composed of two components, the Drell-Yan (i.e. prompt) and the open charm (off vertex) events. To evaluate each contribution, their shapes were obtained from Pythia (details are given in 8); then, the offset distribution was fitted as a superposition of prompt and off-vertex contributions. The fit parameters are the coefficients by which each contribution should be scaled in order to describe the data.

The fit fails to reproduce the data if the prompt yield is forced to 1.1 the expected Drell-Yan yield, and the open charm is left free. In the left panel of fig.2 both contributions are left free: in this way, an accurate representation of data is achieved. The fit parameters indicate that the prompt contribution is two times larger than expected: we can then conclude that the excess is due to a prompt source.

![Figure 2: Left panel: fit to the weighted offset distribution where both the prompt and the open charm yields are free parameters (1.16 < $m_{\mu\mu} < 2.56$). Right: mass spectra of Drell-Yan and excess. The effective temperature of the excess fitted in the $0 < p_T < 2.5$ range for the mass bins of 1.16 - 1.4, 1.4 - 2.0 and 2.0 - 2.56 GeV/c^2 is also shown.](image)

The naive hypothesis that the excess could be due to an increased Drell-Yan yield is ruled out by the $p_T$ dependence of the ratio between the excess and the Pythia-generated Drell-Yan (after acceptance correction). It ranges from 3 at low $p_T$, to 0.5 at high $p_T$, suggesting that the excess is qualitatively different from Drell-Yan. This difference is confirmed by the comparison of the excess and Drell-Yan mass spectra, shown in the right panel of fig.2. The temperatures shown in the figure are obtained from fits in the range $0 < p_T < 2.5$ GeV/c: the systematic error on the temperature determination has been evaluated from fits performed in different $p_T$ ranges and turns out to be of the order of 10 MeV (see also 8).

**J/\psi** suppression in In-In. To measure the amount of anomalous suppression, the $J/\psi$ yield has to be compared to a reference process. Traditionally, the Drell-Yan yield above 4 GeV/c^2 has been used, but this choice increases the statistical error due to the DY limited statistics. To overcome this limitation and fully exploit the $J/\psi$ sample, NA60 has calculated the reference spectrum expected in case of normal suppression only: the relative normalization between the calculated and measured spectra is set to the value resulting from the $\sigma_{J/\psi}/\sigma_{DY}$
analysis (see 9). The resulting suppression pattern is shown in fig.3 (left), compared to the results published by NA60 in Pb–Pb collisions: both sets of data depart from the normal nuclear absorption line for 50 < N<sub>part</sub> < 100, suggesting that N<sub>part</sub> could be well suited as scaling variable between different colliding systems. The data have been compared to theoretical predictions tuned for NA60 In–In collisions: none of them was able to reproduce data, even if the magnitude of suppression is reasonably reproduced (further details in 9).

![Image](image.png)

Figure 3: Left: comparison between the In-In (NA60: square points) and Pb-Pb (NA60: triangle points) suppression patterns. Right: Compilation of the \( \sigma_{J/\psi}/\sigma_{DY} \) values measured in p–A and A-A collisions at the SPS, rescaled, when necessary, to 158 GeV/c. The lines indicate the results of a Glauker fit to the p–A data and the size of the error. The full circle indicates the preliminary NA60 result for p–A collisions at 158 GeV/c.

To evaluate correctly the anomalous suppression, the normal nuclear suppression (which can be extracted from p–A data) must be accurately known. However, up to 2004, p–A data were available only for proton energies of 400 and 450 GeV/c: this implied that the p–A results had to be rescaled to the energy of nuclei beams (158 AGeV/c), under the assumption that \( \sigma_{abs} \) would not change with the beam energy. In order to eliminate uncertainties associated with this assumption, NA60 has taken p–A data at the same energy of the Indium beam on a variety of targets (Be, Al, Cu, In, W, Pb and U). Up to now, a preliminary estimate of \( \sigma_{J/\psi}/\sigma_{DY} \) has been obtained by averaging over the different targets. In fig. 3 right, it is shown a comparison of this result to the previous measurements (rescaled when necessary) as a function of the mean length of nuclear matter traversed by the J/ψ. The NA60 average point corresponds to a length \( \langle L \rangle = 3.4 \text{ fm} \) and falls along the interpolating band, corroborating the correctness of the rescaling procedure and confirming the anomaly of the J/ψ suppression with respect to a pure nuclear absorption scenario.

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CORRELATIONS WITH PHOTONS IN HEAVY-ION COLLISIONS

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We present a study of two-particle correlation functions involving photons and neutral pions in proton-proton and lead-lead collisions at the LHC energy. The aim is to use these correlation functions to quantify the effects of the medium on the jet decay properties.

1 Introduction

Electromagnetic probes have long been thought to be useful to detect the formation of quark-gluon plasma in ultrarelativistic heavy-ion collisions \cite{1,2} and many observables involving photons can be used, since the photon can in principle be used to tag the energy of the recoiling jet. So, by comparing pp and PbPb correlations, one hopes to learn about medium effects in heavy ions collisions. Then, observing a particle recoiling from the photon gives information on the fragmentation properties of this recoiling jet.

2 Model

At leading order (LO) of QCD, the basic two-particle cross section from which we can construct various observables can be written as \cite{3}:

\[
\frac{d\sigma^{AB\to CD}}{dp_{T3} dy_{3} dz_{3} dp_{T4} dy_{4} dz_{4}} = \frac{1}{8\pi s^2} \sum_{a,b,c,d} D_{C/c}(z_3, M_F) D_{D/d}(z_4, M_F) \frac{1}{z_3} \frac{1}{z_4} k_{T3} \delta(k_{T3} - k_{T4}) F_{a/A}(x_1, M) F_{b/B}(x_2, M) \langle M_{12b\to cd}^2 \rangle
\]  

(1)

The medium produced in heavy-ion collisions affects this cross-section by two principal effects, namely: the initial state effects and the final state effects.
2.1 Initial state effects

These effects result from the modification of the structure functions by the shadowing and anti-shadowing effects which are hard to calculate theoretically. So, we use here the parametrization of Eskola et al., 4, who tabulate a function $S_{a/A}(x, M)$ in which quarks and gluons are treated separately and relates the PDFs in nucleon $N$ to those in a nucleon $A$ via

$$ F_{a/A}(x, M) = S_{a/A}(x, M) F_{a/N}(x, M) $$

(2)

Nuclear effects in PDFs produce small changes in high-$p_t$ particle production at RHIC and LHC, ”at most 25% “.

2.2 Final state effects

In a medium, quarks and gluons lose energy by radiating gluons. Their fragmentation is modified (See BDMPS). 5, 6. One can define a medium-modified fragmentation function (FF) $D_{D/\ell}^{med}(z_d, M_F, k_{Td})$ which is calculated from the medium-induced BDMPS gluon spectra $dI/d\omega_d$.

$$ z_d D_{D/\ell}^{med}(z_d, M_F, k_{Td}) = \int_0^{k_{Td}(1-z_d)} d\epsilon D_{D/\ell}(\epsilon, k_{Td}) z_d^\star D_{D/\ell}(z_d^\star, M_F) $$

(3)

where $z_d = \frac{\epsilon + \omega_d}{k_{Td}}$ and $z_d^\star = \frac{\epsilon + \omega_d}{\epsilon + k_{Td}}$. The BDMPS energy loss distribution is characterized by the energy scale $\omega_d$.

$$ \omega_d = \frac{1}{2} \hat{q} L^2 $$

(4)

The so-called gluon transport coefficient $\hat{q}$ reflects the medium gluon density and $L$ is the length of matter covered by the hard parton in the medium.

2.3 Medium-modified fragmentation functions

![Graphs showing medium-modified FFs](image)

Figure 1: (left): Medium-modified FFs $D_{D/\ell}^{med}(z_d, M_F, k_{Td})$ for various energy loss scales, $\omega_d = 0$, 25 and 50 GeV. The parton energy is $k_{Td} = 50$ GeV and (right): Ratio of medium-modified ($\omega_d = 50$ GeV) over vacuum ($\omega_d = 0$ GeV) FFs for various parton energy, $k_{Td} = 25, 50$ and 100 GeV. The fragmentation scale is set to $M_F = p_t/2$

We observe that effects of parton energy loss become more pronounced as $z$ gets larger, due to the restricted available phase space in Eq. (1). Gluons lose more energy than quarks do from their larger color charge ($C_g = 3, C_q = 4/3$). In the high energy limit $k_{Td} \gg \omega_d$ and thus $z^\star \simeq z$, the medium effects vanish and the ratio approaches one.
3 The correlations

We construct from Eq. (1) the following observables:
- the invariant mass of the particle pair, $m_{34}$
- the transverse momentum of the pair, $q_T = |\vec{p}_{T3} + \vec{p}_{T4}| = k_T |z_3 - z_4|
- the relative transverse momentum of the particles (also called momentum balance), $z_{34} = \frac{\vec{p}_{T3} \cdot \vec{p}_{T4}}{p_{T3} p_{T4}} = \frac{z_3}{z_4}$

Note that for direct photon, momentum fraction is $z_3 = 1$. For fragmentation photon, $z_3 < 1$ and is further reduced by medium-induced energy loss.

3.1 Phenomenology of $\gamma - \pi^0$ correlations

The photon can be produced directly and the recoiling jet fragments into a pion (labeled 1f), or both the photon and the pion are produced by fragmentation of partons (labeled 2f).

For LHC, we impose the following cuts: $p_{T\gamma} \geq 5$ GeV and $p_{T\pi} \geq 25$ GeV. We observe clearly in Fig.2 the expected effect of the suppression. Energy loss effects modify the distributions much more drastically.

3.2 Phenomenology of $\gamma - \gamma$ correlations

Figure 2: (left): The four distributions in $\gamma - \pi^0$ production for p-p (open dots) and Pb-Pb scattering ($\omega_c = 0$ GeV; black squares; $\omega_c = 50$ GeV; open squares) at $\sqrt{s} = 5.5$ TeV, $|\eta| < 0.5$, $|\eta| < 0.5$; cuts imposed are $p_{T\gamma} > 25$ GeV, $p_{T\pi} > 5$ GeV. (right): The same as left Figure but the distributions are normalized to p-p case.

Figure 3: (left): Top: The $p_{T\gamma}$ and $q_T$ transverse momentum distributions in $\gamma - \gamma$ production for p-p (open dots) and Pb-Pb scattering ($\omega_c = 0$ GeV; black squares; $\omega_c = 50$ GeV; open squares) at $\sqrt{s} = 5.5$ TeV, with the same cuts as in Figure (2). Bottom: The same distributions normalized to the p-p case. (right): The $z_{34}$ distribution in $\gamma - \gamma$ for p-p (open dots) and Pb-Pb scattering (for the same kinematical cuts). Bottom: The same distributions normalized to the p-p case.
The new feature here is that both photons can be produced directly (direct process) in which case they are not affected by medium. The study of the $\gamma - \gamma$ correlations is made in the same kinematic regime as before, i.e. $p_{T,\gamma} > 25$ GeV and $p_{T,\gamma} > 5$ GeV. In particular, we observe in Fig.3 (right: Bottom) a strong suppression as $z_{34}$ gets close to 1, as expected from the restricted phase space in Eq. (3). Assuming the 1-fragmentation dominates, we have at LO $z_{34} \simeq z$. The distribution $dz/dz_{34}$ is thus reminiscent of the photon fragmentation function, $D_{\gamma/jk}(z_{34})$. This observable offers therefore a "direct" access to the medium-modified fragmentation functions. On the same Figure (left panel), the photon $p_T$ spectra are also determined (lower left), the quenching is maximal around the upper cut. Indeed, as $p_T$ approaches the upper cut from "below", events with larger $z$ are selected, $p_{T1} \simeq p_{T2}$, where energy loss effects are most pronounced.

The $z$ distributions in Fig. 4 appear to follow closely the input functions in Fig. 1 and therefore provide

![Figure 4: Example of a photon-photon correlation function at LHC^a.](image)

a way to probe in detail the jet energy loss mechanism. The same measure involving a pion (photon-pion correlation) has a larger rate but leads to a more complex picture because of the convolution with the production processes: 1f at small $z_{34}$ and 2f at large $z_{34}$.

**Conclusions**

We have discussed various photon tagged correlations as a tool to study jet fragmentation in hot medium created in heavy-ion collisions. We show that significant effects could be expected at LHC energy both in the $\gamma - \pi^0$ and $\gamma - \gamma$ channel. The use of asymmetric cuts in the transverse momentum of both particles allow the possibility to map out the parton fragmentation functions modified by the medium. The variety of observables presented here should help constrain the underlying model for parton energy loss. We believe our present LO predictions to be valid roughly up to $z_{34} \simeq 0.8$.

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Experimental tests of small-$x$ QCD

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Current and future experimental studies of the high-energy limit of QCD, dominated by non-linear gluon dynamics in the low-$x$ sector of the hadron wavefunctions, are presented. Results at HERA (proton) and RHIC (nucleus) pointing to the possible onset of parton saturation phenomena, and perspectives at the LHC and new proposed DIS facilities are outlined.

Keywords: QCD, low-$x$, gluon saturation, HERA, RHIC, LHC.

1 Introduction

Gluons provide the dominant contribution to hadronic scattering cross sections at high-energies. Deep-Inelastic (DIS) experiments of electrons off protons at the HERA collider at DESY have shown that for values of the parton momentum fraction $x = p_{\text{parton}}/p_{\text{proton}} \lesssim 0.01$, the proton wavefunction is basically purely gluonic (Fig. 1). This is so, because gluons are “cheap” to radiate: the probability of emitting a gluon increases as $\propto \alpha_s \ln(Q^2)$ and $\propto \alpha_s \ln(1/x)$ according to the standard linear QCD evolution (DGLAP\(^1\) and BFKL\(^2\) resp.) equations. As a matter of fact, for decreasing values of $x$ the gluon density increases so fast that unitarity would be ultimately violated, even for processes with large virtualities $Q^2 \gg \Lambda^2_{\text{QCD}}$. The theoretical expectation\(^3,4\) is that at some small enough value of $x$ ($\alpha_s \ln(1/x) \gg 1$) non-linear gluon-gluon fusion effects – not accounted for in the DGLAP/BFKL equations – will become important and will tame the growth of the parton densities. The onset of saturation in the proton (or in a nucleus with $A$ nucleons) is expected for parton momenta $Q^2 \lesssim Q_s^2$ where $Q_s$ is a dynamical “saturation scale”\(^3,4\) which depends on the transverse size ($\pi R^2$) of the hadron:

$$Q_s^2(x) \simeq \alpha_s \frac{1}{\pi R^2} \frac{1}{x} G(x, Q^2) \sim A^{1/3} x^{-\lambda} \sim A^{1/3} (\sqrt{s})^\lambda \sim A^{1/3} e^{\lambda y},$$

Eq. (1) tell us that $Q_s$ grows with the energy of the collision, $\sqrt{s}$, or equivalently, with the rapidity of the parton $y = \ln(1/x)$. The nucleon number $A$ dependence implies that, at equivalent energies, saturation effects will be amplified by factors as large as $A^{1/3} \approx 6$ in heavy nuclear targets ($A = 208$ for Pb) compared to protons. The regime of high gluon densities is often described in terms of the colour-dipole\(^6,7\) or “Colour Glass Condensate” (CGC)\(^8\) effective theories, with the corresponding non-linear BK/JIMWLK\(^9,10\) evolution equations.

2 Gluon saturation at HERA?

Although the arguments for saturation are well justified theoretically, no strong deviation from the linear QCD equations has been conclusively observed in the perturbative kinematical range
covered at HERA. Most of the experimental observables, in particular those of more inclusive nature such as the total $\sigma_{\gamma p}$ cross-section $d^2\sigma/dx dQ^2 \approx 2\pi \alpha_s^2/(xQ^4)F_2(x,Q^2)$, are in good accord with the standard DGLAP expectations (Fig. 1). However, it is worth to note that the saturation scale at HERA energies is in a regime of relatively low virtualities, $Q_s^2 \approx 1$ GeV$^2$ and thus – since Bjorken $x$ and virtuality are correlated as $x \approx Q^2/s$ – the interpretation of most of the truly low-$x$ range probed ($x \lesssim 10^{-4}$) is “blurred” by its proximity to the non-perturbative regime. As a consequence, the gluon distribution function $xg(x,Q^2)$ indirectly obtained from the $F_2$ scaling violations – via $\partial F_2(x,Q^2)/\partial \ln(Q^2) \approx 10 \alpha_s(Q^2)/(2\pi x)g(x,Q^2)$ – is poorly constrained below $x \approx 10^{-4}$. Different DGLAP parametrizations based on DIS-only data (ZEUS-PDF, H1-PDF, Alekhin02) or on global fits (CTEQ6.5M, MRST-NLO) yield gluon PDFs differing by factors of 3 or more (Fig. 1, right).

Notwithstanding those uncertainties, there are three empirical observations at HERA (summarized in Fig. 2) that favour a possible onset of low-$x$ parton saturation in the proton. The leftmost plot shows the geometric scaling\textsuperscript{7} property of inclusive $\sigma_{\text{DIS}}$ which, instead of being a function of $x$ and $Q^2$ separately, for $x < 0.01$ it features a single dependence on the parameter $\tau = Q^2/Q_s^2(x)$ where $Q_s^2(x) = Q_0^2(x_0)^4$ with $\lambda \sim 0.3$, $Q_0 = 1$ GeV, and $x_0 \sim 3 \times 10^{-4}$. Such a scaling property is naturally explained by gluon saturation models, whereas the DGLAP approach can only reproduce it via a fine tuning of the initial parameterization of the gluon distribution used. Another piece of evidence for saturation effects at HERA is provided by diffractive processes, where the proton remains intact after the “quasi-elastic” interaction with the photon. Diffractive scattering, accounting for 10–15% of the total DIS cross-section, is characterized by colourless two-gluon exchange, and thus it constitutes a sensitive probe of the gluon densities\textsuperscript{11}. Surprisingly, the ratio of the diffractive to total $\sigma_{\gamma p}$ cross-sections is found\textsuperscript{12} to be roughly constant as a function of the center-of-mass energy $W$ and $Q^2$ (Fig. 2, center). This is in disagreement with the naive pQCD expectations of an increase of the ratio according to $r_{\text{diff}}/\sigma_{\text{tot}} \sim 1/\sqrt{xg(x,Q^2)}^2 \sim W^{4\lambda}/W^{2\lambda} \sim W^{2\lambda}$. The last indication of a

\textsuperscript{a}See http://durpdg.dur.ac.uk/hepdata/pdf3.html
“tension” between the standard linear QCD equations and low-$x$ HERA data comes from the
longitudinal structure function, $F_L(x, Q^2)$, which at variance with $F_2$, is directly proportional
to $x g(x, Q^2)$. The $F_L(x, Q^2)$ derived from DGLAP analyses becomes unphysically negative below $x \approx 10^{-4}$ for relatively small $Q^2$ values, whereas it is a well-behaved object in saturation
models (Fig. 2, right).

![Graphs showing saturation effects at HERA](image)

Figure 2: Hints of saturation at HERA. Left: Geometric scaling $^7$ of $\sigma_{et}^p$ as a function of $\tau = Q^2/Q_0^2$ for $x < 0.01$. Center: Ratio of diffractive to total DIS cross-sections $^{12}$ as a function of $W$ and $Q^2$. Right: Longitudinal structure
function, $F_L(x, Q^2 = 2 \text{ GeV}^2)$, predicted by DGLAP $^{13}$ and saturation $^{15,13,12}$ models.

3 Gluon saturation at RHIC?

Among the interesting observations in nucleus-nucleus (A-A) collisions at RHIC is the possible
onset of parton saturation phenomena. Though nuclei at RHIC are probed at lower energies
($\sqrt{s_{NN}} = 200 \text{ GeV}$) than protons at HERA, saturation effects are “amplified” thanks to the
increased transverse parton density in the former compared to the latter. Two empirical observations support the Color-Glass-Condensate (CGC) predictions of a reduced parton flux in the
incoming ions due to enhanced non-linear QCD effects. On the one hand, the measured hadron multiplicities $^{16,17,18,19}$, $dN_{ch}/d\eta|_{\eta=0} \approx 700$, are significantly lower than the $dN_{ch}/d\eta|_{\eta=0} \approx 1000$ values predicted by minijet $^{20}$ or Regge $^{21}$ models, but are well reproduced by CGC
approaches $^{22}$. Parton multiplicity distributions at high energies are perturbatively calculable in
saturation approaches, since they are governed by a semi-hard saturation scale $Q_s^2 \propto \sqrt{s_{NN}}^{-\lambda}$
with an exponent $\lambda$ constrained by e-A data $^{23}$. Assuming parton-hadron duality, hadron multiplicities at mid-rapidity rise then proportionally to $Q_s^2$ times the transverse (overlap) area, a feature that accounts naturally for the experimentally observed factorization of $\sqrt{s_{NN}}$- and centrality-dependences in $dN_{ch}/d\eta|_{\eta=0}$ (Fig. 3, left).

The second manifestation of CGC-like effects in the RHIC data is the BRAHMS observation $^{19}$ of
suppressed yields of semi-hard hadrons ($p_T \approx 2 - 4 \text{ GeV}/c$) in d-Au relative to p-p collisions
at increasingly forward rapidities (up to $\eta \approx 3.2$, Fig. 3, right). Hadron production at such small
angles is theoretically sensitive to partons in the Au nucleus with $x_{min}^{Au} = (p_T/\sqrt{s_{NN}}) \exp(-\eta) \approx O(10^{-3})$ $^{24}$. The observed nuclear modification factor, $R_{dAu} \approx 0.8$, cannot be reproduced by pQCD calculations $^{24,25,26}$ that include the same leading-twist nuclear shadowing that describes the d-Au data at $y = 0$, but can be described by CGC approaches that parametrise the Au
nucleus as a saturated gluon wavefunction $^{27,28}$.
4 Low-\(x\) QCD studies at the LHC

The Large Hadron Collider (LHC) at CERN will provide p-p, p-A and A-A collisions at \(\sqrt{s_{NN}} = 14, 8.8\) and 5.5 TeV with luminosities \(L \sim 10^{34}, 10^{29}\) and \(5 \times 10^{30}\) cm\(^{-2}\) s\(^{-1}\) respectively. Following Eq. (1), the relevance of low-\(x\) QCD effects will be significantly enhanced due to the increased: center-of-mass energy, nuclear radius (\(A^{1/3}\)), and rapidity of the produced partons\(^{29,30}\). At the LHC, the saturation momentum \(Q^2_s \approx 1\) GeV\(^2\) (proton) – 5 GeV\(^2\) (Pb) will be more clearly in the perturbative regime\(^5\), hard probes will be copiously produced, and the \(x\) values experimentally accessible will be much lower than at previous colliders: \(x \approx 10^{-3}\) (10\(^{-6}\)) at central (very forward) rapidities. All LHC experiments have interesting detection capabilities in the forward direction which will help to constrain the PDFs in the very low-\(x\) regime: (i) CMS\(^{31,32}\) can measure inclusive jet and Drell-Yan production down to \(x \sim 10^{-6}\) using the CASTOR calorimeter at rapidities \(5.5 < \eta < 6.6\) as well as Mueller-Navelet dijets (very sensitive to non-DGLAP evolution)\(^{33,34}\) separated by rapidities as large as \(\Delta \eta \sim 10\); (ii) ALICE and LHCb feature a forward muon spectrometer (covering \(2 \lesssim \eta \lesssim 5\) which gives them access to heavy-quark, quarkonia and gauge boson measurements\(^{35}\) down to \(x \sim 10^{-5}\).

The advance in the study of low-\(x\) QCD phenomena will be especially substantial in Pb-Pb collisions. In saturation models, there is a one-to-one correspondence between the effects of rapidity- and \(\sqrt{s_{NN}}\)-dependences, because a parton distribution boosted to higher rapidity \(\eta\) is equivalent to a distribution sampled in a process at higher \(\sqrt{s_{NN}}\). As a consequence, the saturation physics explanation of the increasing forward suppression of semi-hard yields at RHIC will imply a significant A-A hadron suppression at LHC mid-rapidities. CGC predictions for charged hadron multiplicities in central Pb-Pb at 5.5 TeV\(^{36}\) are \(dN_{ch}/d\eta|_{\eta=0} \approx 1500\), i.e. 3-4 times lower than the pre-RHIC era results. [As a matter of fact, such low multiplicities help CMS\(^{37}\) and ATLAS\(^{38}\) to become very competitive experiments in the heavy-ion running mode]. Arguably, one of the cleanest way to study the low-\(x\) structure of the Pb nucleus at the LHC is via Ultra Peripheral collisions (UPCs)\(^{39}\) in which the strong electromagnetic fields (equivalent flux of quasi-real photons) generated by the colliding nuclei can be used for photoproduction studies at maximum energies \(\sqrt{s_{NN}} \approx 1\) TeV, 3-4 times larger than at HERA. Full simulation+reconstruction studies\(^{37,32}\) of quarkonia photoproduction (\(\gamma\) Pb \(\rightarrow\) J/\(\Psi\)) tagged...
with very-forward neutrons, show that CMS can carry out detailed $p_T, \eta$ measurements in the dielectron and dimuon $\Upsilon$ decay channels. Such processes probe $x$ values in the nucleus as low as $x \sim 10^{-4}$ (Fig. 4, left).

![Kinematic $(x, Q^2)$ plane probed in $e, \gamma, A$ processes: existing data compared to (a) UPC $\Upsilon$ photoproduction processes (left), and (b) new proposed nuclear DIS facilities (right): LHeC and EIC/eRHIC.](image)

Figure 4: Kinematic $(x, Q^2)$ plane probed in $e, \gamma, A$ processes: existing data compared to (a) UPC $\Upsilon$ photoproduction processes (left), and (b) new proposed nuclear DIS facilities (right): LHeC and EIC/eRHIC.

5 Proposed future deep-inelastic facilities

Two different collider projects – the Large Hadron Electron Collider (LHeC) $^{40}$ and the Electron Ion Collider (EIC) $^{41}$ – have been recently proposed to study deep-inelastic lepton-hadron (e-p, e-d and e-A) scattering for momentum transfers $Q^2$ as large as $10^6$ GeV$^2$ and for Bjorken $x$ down to the $10^{-6}$. Both projects have a strong focus on the study of the low-$x$ gluon structure of protons and nuclei and on non-linear QCD evolution. LHeC (EIC/eRHIC) proposes to add an extra $E_e = 70$-GeV (100-GeV) electron ring to the LHC (RHIC/JLab) proton/nucleus collider(s).

In the nuclear DIS sector, LHeC (EIC/eRHIC) would allow e-A collisions at $\sqrt{s_{NN}} = 2\sqrt{s_N E_e} = 880$ (63) GeV. Fig. 4 right, shows the $(x, Q^2)$ ranges accessible to both machines compared to existing nuclear DIS data. Both proposed facilities would significantly extend the (meager) kinematical regime of current data, fully mapping out the range of Bjorken-$x$ at virtualities around the saturation scale ($Q^2_s \approx 1 - 10$ GeV$^2$, indicated by the black curve in the plot) and providing very valuable insights on the high-energy limit of QCD.

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Exploring the physics of strong color fields with a new Electron-Ion Collider

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The proposed polarized electron-ion collider (EIC) will allow for precision exploration of various novel aspects of QCD including low-\(x\) phenomena, the spin structure of the proton, and gluon saturation in heavy nuclei. As this project gains momentum, it is increasingly important for the QCD community to understand quantitatively the kinematical reach and expected sensitivities for various measurements. We briefly summarize key accelerator design parameters and then focus on expected measurement sensitivities, thus exposing how the EIC will allow an extension of the successful HERA program into exciting new regimes.

Keywords: select here a few keywords for your article, from the A&A thesaurus

QCD is a spectacularly successful theory, yet it remains an active field of research in particle physics. While it has withstood over two decades of tests, we have yet to understand fully the mechanisms by which complex and rich phenomena emerge from a theory based solely on symmetry and local gauge invariance.\(^1\) This quest is being pursued on many fronts: parton distributions, diffraction, shadowing, hadron spin structure, low-\(x\) saturation, and relativistic heavy ion collisions, to name a few. There is little doubt that our current understanding of QCD, in particular the role of the gluon in QCD dynamics, has been greatly advanced by using electron-proton collisions at the HERA collider, which was the first facility to probe with high luminosity deep into the regime where gluons play a dominant role in hadron structure. In the following, we discuss how a high luminosity Electron Ion Collider (EIC) will enable similar advances in the study of strong color fields, in particular many body gluon states probed at the saturation scale. The crucial element in these studies will be the use of heavy nuclei that, when probed at small Bjorken \(x\), amplify such novel gluon states.

From deep inelastic scattering (DIS) we know that gluons carry half of the momentum of a proton, and from HERA we know that gluons dominate for \(x < 0.01\).\(^2\) However, when probed at sufficiently small scales and at low-\(x\), it is predicted that the gluon distribution within a proton will saturate. There are many quantitative arguments for this, but it is rather intuitive that at sufficiently high gluon density \(2 \rightarrow 1\) gluon fusion (non-linear dynamics) will begin to dominate over \(1 \rightarrow 2\) gluon splitting (linear dynamics). This transition occurs at a scale commonly referred to as the saturation scale \((Q_s)\). Geometrical considerations\(^3\) show that, in nuclei, \(Q_s \propto A^{1/3}x^{-\delta}\), where \(\delta \sim 1/3\).\(^4\) Thus, heavy nuclei can be used to amplify the scale at which gluon saturation sets in. Figure 1 shows the saturation scale in the \(Q^2\) vs \(x\) plane for ions ranging from protons to Au.\(^2\) Additionally, the colored diagonal lines show the kinematic limits of the EIC for various beam energies. The shaded band illustrates the kinematically accessible saturation region for 20 GeV electrons colliding with 100 GeV Au nuclei. The saturation scale for gluons in Au can be accessed at a factor of 10 larger \(x\) and \(Q^2\) than \(Q_s\) for gluons in a proton. This has two
significant impacts. First, by substituting heavy ions for protons, one can access the saturation regime at lower beam energies—a dramatic savings in cost. Second, $Q_{s,g}^2(Au) > A_{QCD}$, thereby enabling use of perturbative methods for calculation. The saturation regime can be reached with reasonably modest beam energies, and it can be explored with the use of perturbative calculations.

In DIS the differential cross section $\frac{d^2 \sigma}{dx dq^2}$ can be decomposed in terms of two structure functions ($F_2(x, Q^2)$ and $F_L(x, Q^2)$), where $F_2$ is directly sensitive to (anti)quark distributions, and $F_L$ is directly proportional to the gluon momentum distribution. These structure functions provide a direct means to quantitatively study saturation phenomena and we discuss them in more detail below. In the following we show the measurement prospects for various EIC configurations with a maximum center-of-mass energy of $\sqrt{s} = 14 \sim 140$ GeV/n and a maximum luminosity $L = 10^{33} \sim 10^{35}$ cm$^{-2}$s$^{-1}$, a factor of 100 higher than HERA. As the world’s first electron heavy ion collider, the EIC would enable the high precision exploration of $F_2$ and $F_L$ of heavy Nuclei in the saturation regime, truly term $\text{incom} \text{nuita}$. As we will show, the luminosity, energy and collider kinematics will be used to differentiate competing models of low-$x$ QCD phenomena.

One of the first measurements at the EIC will be $F_2(x, Q^2)$ for both heavy (Au) and light (d) ions. The ratio, shown vs $Q^2$ for four $x$ bins, is shown in Figure 2. The points represent the anticipated statistical precision achievable with an integrated luminosity of 4/A fb$^{-1}$. The colored lines are predictions from competing models. $F_2$ is directly sensitive to quark distributions, and is sensitive to gluons via scaling violations. nDS, EKS and FGS are pQCD models with differing treatment of parton shadowing, and they are compared to predictions from the Color Glass Condensate (CGC) model. Within the expected precision of the measurements, differentiation between the different models is clearly possible in the region $10^{-4} < x < 10^{-2}$.

With the ability to accelerate both light and heavy ions over a wide range of energies, the EIC will be able to make significant contributions to the understanding of the gluon distribution in the proton. At low $x$, $F_L(x, Q^2) \propto \alpha_s x \ G(x, Q^2)$, where $G(x, Q^2)$ is the gluon distribution. Extraction of $F_L$ requires running at multiple beam energies, a task that highlights the flexibility of the EIC. Figure 3 shows $F_L(x)$ for a proton. The red points are from existing NMC fixed
Figure 2: The ratio of $F_2(Q^2)$ in Au to that in d nuclei, for four bins in $x$. The symbols represent the anticipated statistical precision achievable. The curves represent models with differing treatment of low-$x$ phenomena. Color online.

Figure 3: $F_L(x)$ for protons from fixed target (red), projected H1 (blue) and projected EIC (black). Color online.
target data, and the blue H1 points show the expected precision (statistical and systematic uncertainties) achievable from the recent HERA energy scan. The black points show the achievable precision from one year of running the EIC at four different energies (statistical uncertainties only). The EIC measurements will clearly compliment the HERA results, as well as bridge the region between HERA and fixed target results, contributing to a direct measurement of the gluon distribution for \(10^{-4} < x < 10^{-1}\).

Finally, Figure 4 shows the ratio of gluon distributions in Pb to that in d \(\left(\frac{R^{p}}{R^{d}} = \frac{F_{L}^{p}}{F_{L}^{d}} \sim \frac{G_{p}}{G_{d}}\right)^{2}\). Current data on the gluon distribution at low-\(x\) in heavy nuclei is sparse. In turn, constraints on theoretical models are weak, as shown by the vast range of different theoretical predictions. The gluon distribution is critical for accurate calculations of cross sections at both RHIC and the LHC. With its high luminosity, the EIC can make significant contributions to the understanding of the gluon distribution in heavy nuclei over the relevant \(x\) range.

In conclusion, as the world’s first high energy electron (heavy) ion collider, and with a luminosity approximately one hundred times that of HERA, the EIC will allow precision exploration of strong color fields. The use of heavy nuclei will amplify the scale at which saturation phenomena are predicted, placing it well within the accessibility of the EIC. There are many topics we have neglected to discuss, in particular diffraction, spin decomposition of the proton, and the study of partonic energy loss in cold nuclear matter. The physics program of the EIC is rich, diverse, and well targeted toward a unified understanding of strongly interacting matter. The project is gaining momentum on an international scale and will provide a continuation of the successful HERA program into exciting new regimes of QCD.

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References

HEAVY ION PHYSICS WITH THE ALICE EXPERIMENT AT LHC

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ALICE is the experiment at the LHC collision at CERN dedicated to heavy ion physics. In this report, the ALICE detector will be presented, together with its expected performance as far as some selected physics topics are concerned.

1 Introduction

Besides its proton-proton (pp) physics program, the Large Hadron Collider at CERN, which is going to start operating in early 2008, will deliver lead-lead (PbPb) collisions at a centre-of-mass energy \( \sqrt{s_{NN}} = 5.5 \text{ TeV} \), opening the door to unexplored regimes in heavy-ion particle physics from several points of view, such as the centre-of-mass energy, the energy density reached, the lifetime and the volume of the Quark-Gluon Plasma (QGP) system which may be created in the collision. ALICE (A Large Ion Collider Experiment) will be the experiment dedicated to study heavy-ion collisions at the LHC. Its physics program will span over a large number of observables, from the global properties of the collisions (multiplicities, rapidity distributions...), to more selective QGP signals (like direct photons, chamonium and bottomonium...). In this report, an overview of the LHC and the ALICE detector will be first provided. Then, the physics performance of the experiment as far as open charm and open beauty detection, quarkonia and jet physics studies will be discussed. Finally, the ALICE “First Day” Physics will be briefly presented.

2 The ALICE Experiment at the LHC

The Large Hadron Collider experimental program will deal with different collision systems. Apart from the main pp and PbPb runs, also asymmetric collisions (such as p-Pb) and collisions with heavy ions other than Pb (e.g. Ar) will be studied. Table 1 quotes the expected values for the most significant running parameters for the pp and PbPb collision systems, which have been used as a reference for the results presented hereafter.

ALICE will be the experiment at the LHC dedicated to heavy-ion physics. In the central rapidity region \(-0.9 < \eta < 0.9\), inside the L3 magnet (providing the experiment with a weak solenoidal magnetic field \( B = 0.2 - 0.5 \text{ T} \)), ALICE will be endowed with subdetectors specialized in tracking and identifying the particles produced in the collisions, namely the ITS (Inner Tracking System), the TPC (Time Projection Chamber), the Transition Radiation Detector (TRD), and the Time Of Flight (TOF). The central region will be also equipped with two detectors having partial azimuthal coverage: the HMPID (High Momentum Particle Identification Detector) and the PHOS (Photon Spectrometer). Besides, an Electromagnetic Calorimeter (EMCAL) has recently been added to the ALICE central setup, covering the pseudorapidity range \( |\eta| < 0.7 \) and \( 110^\circ \) in azimuth. To be noted that, due to the late inclusion of the EMCAL in the ALICE design, its installation will be slightly delayed with respect to the other central detectors. In the pseudorapidity region \(-4.0 < \eta < -2.5\), instead, the ALICE Muon Spectrometer will be the detector enchanced of studying both heavy quark vector mesons and the \( \phi \) meson via the \( \mu^+\mu^- \) decay channel, and the production of open heavy flavours. At large rapidity values, other detectors will be installed, namely the Zero Degree Calorimeter (ZDC), the Photon Multiplicity Detector (PMD), the Forward Multiplicity Detector (FMD) and the V0 and T0 detectors. Finally, the ALICE triggering on cosmic rays will be performed by the ACORDE detector (for more details see ).
Table 1: LHC expected running conditions for pp and PbPb collisions. Due to the limited ALICE rate capability, in the experiment interaction region a lower pp luminosity value is foreseen (quoted within brackets in the first row, third column of the Table).

<table>
<thead>
<tr>
<th>Collision System</th>
<th>$\sqrt{s_{NN}}$ (TeV)</th>
<th>$L_0$ (cm$^{-2}$.s$^{-1}$)</th>
<th>Run Time (s/year)</th>
<th>$\sigma_{wom}$ (b)</th>
</tr>
</thead>
<tbody>
<tr>
<td>pp</td>
<td>14.0</td>
<td>$10^{24}$ (10$^{31}$)</td>
<td>10$^7$</td>
<td>0.07</td>
</tr>
<tr>
<td>PbPb</td>
<td>5.5</td>
<td>$10^{27}$</td>
<td>10$^6$</td>
<td>7.7</td>
</tr>
</tbody>
</table>

3 ALICE Physics Performance

The ALICE experiment will study a wide number of observables, which will require the use of various experimental techniques. In the following, the ALICE expected performance in terms of measurement of heavy flavour (Sect. 3.1), quarkonia (Sect. 3.2) and jet production (Sect. 3.3) will be briefly overviewed.

3.1 Heavy Flavour Production in the Central Detectors

Due to their high mass, heavy flavour quarks (charm and beauty) are produced at the very early stage of the collision. For this reason, the measurement of open charm and open beauty production is one of the main observables which can trace the initial phase of the collision and provide information on the possible formation of the Quark-Gluon Plasma. Experimentally, the effects of medium modifications on the final states have shown up in the parton energy loss observed in AA collisions at RHIC, where a departure from binary scaling not only for light charged hadrons, but also for heavy flavour production (for "non-photonic electrons", in fact, which are likely to come from heavy flavour decays) has been observed.

At the LHC, detailed studies on heavy flavour production will be even more feasible compared to RHIC, thanks to the high expected yields for open charm and open beauty, which are expected to be of the order of $\sim 100 \, c\bar{c}$ and $\sim 5 \, b\bar{b}$ in a central PbPb collision. To be noted that these analyses will be important also because they will be used as a reference for quarkonia production studies (see Sec. 3.2).

The open charm and open beauty analysis will rely on the capabilities of the ALICE vertex detector, the ITS, which will be characterized by a resolution on the measurement of the track impact parameter better than 100 $\mu$m for $p_T > 0.6$ GeV/c, allowing to fully reconstruct heavy flavour decays. The tracking and the momentum measurement will be provided by the TPC detector, while particle identification will be performed by the TOF detector in the case of charged hadrons, and by the TRD and the TPC in the case of electrons. As an example of the ALICE performance in measuring open charm, the left panel of figure 1 shows the K$\pi$ invariant mass distribution for the study of the $D \rightarrow K^- \pi^+$ channel, as obtained for $10^7$ PbPb central events.

Being able to measure separately charm and beauty hadrons will offer ALICE the opportunity to study the dependence of the energy loss on the mass of the heavy parton. Besides, the heavy-to-light ratio $R_{D^{(*)}}(p_T)$, defined as the ratio of the nuclear modification factor of D mesons to that of charged light-flavoured hadrons, will allow to shed light on the colour-change dependence of QCD energy loss. For more details on open charm and open beauty analyses with the ALICE detector, see.

3.2 Quarkonia

The study of quarkonia production in heavy ion collision is one of the main observables which can be used to investigate the properties of the medium created in the collision. As a matter of fact, quarkonia

---

*a*The quoted values for the open charm and open beauty production rate have been computed taking into account also the shadowing effect.

*b*The centrality class of an event has been determined according to the impact parameter $b$. In particular, the PbPb 5% most central events correspond to a selection on the event impact parameter $b < 3.5$ fm.
are expected to be sensitive to the collision dynamics (both at short and long timescales) and to plasma formation. Moreover, since different quarkonia dissociate at different temperatures, they can serve as a thermodynamic probe of the medium\(^5\).

From an experimental point of view, the expected anomalous suppression of the \( J/\Psi \) related to Debye colour screening\(^6\) observed at the CERN SPS\(^7\) was thought to be even stronger at RHIC, even in case some dissociation mechanisms due to comoving hadrons would have occurred. Contrarily, the level of \( J/\Psi \) suppression was found to be the same as that at the SPS\(^8\), which would imply the presence of some recombination effect\(^9\) competing with the suppression one. As a consequence, at the LHC, because of the even higher \( \bar{c}c \) rate, the regeneration mechanism is expected to dominate, so that an enhancement in the \( J/\Psi \) yield may even take place\(^10\).

The measurement of quarkonia in ALICE will be performed both at midrapidity in the dielectron channel (making use of the ITS and TPC for tracking, and of the TRD for electron identification), and in the forward rapidity regions, where the ALICE Muon Arm will study the dimuon channel. This implies that the ALICE detector will be able to measure quarkonia production in two complementary Bjorken-x regions, allowing to investigate the PDFs in the nuclei. Moreover, the feed-down from B decays for \( J/\Psi \) production will be kept under control thanks to the open beauty measurements presented in Sec. 3.1.

The right panel of figure 1 shows the ALICE expected performance in a 10% PbPb run (5% most central events) for the measurement of the dimuon spectrum for charmonium (\( J/\Psi \), left) and bottomonium (\( \Upsilon \), right). It has been shown\(^3\) that for the \( J/\Psi \) the statistics collected during one PbPb data taking period will be enough to measure also the \( p_T \) dependence of the charmonium spectrum, while for the \( \Upsilon \) case, because of the smaller statistics, either different centrality classes of events will have to be merged, or two or three data taking periods will be necessary. Finally, it is worth to mention that for quarkonia studies the sensitivity of the results to different suppression scenarios has been investigated. For more details, see\(^3\).

### 3.3 Jet Physics

Proton-proton collisions at the LHC will be characterized by a very high jet production rate. On one hand, jets with energy larger than 20 GeV are expected to occur with a frequency of \( \sim 1 \) per PbPb central event; on the other, \( \sim 10^5 \) highly energetic jets with \( E_T \sim 200 \) GeV are expected to be produced in \( 10^6 \)s of PbPb data taking. Thanks to their high transverse energy, it will be possible to single out
and reconstruct these high energy jets on an event-by-event basis also in a very entangled environment as that produced in a PbPb event. Moreover, the addition of the Electromagnetic Calorimeter detector (EMCAL) in ALICE will improve the performance of the experiment in terms of jet physics, both at the level of energy measurement, and at the level of the high energy jet triggering capabilities of ALICE.

One of the main characteristics of ALICE jet reconstruction is that in order to keep under control as much as possible the background level, it will make use of a limited jet cone size, \( R_c = \sqrt{(\Delta \eta)^2 + (\Delta \phi)^2} \approx 0.3 \div 0.4 \). Moreover, appropriate \( p_T \) cuts will have to be applied, as shown in Figure 2. Here, the energy of the jet from charged particles as a function of the cone size \( R_c \) is drawn for different energies. The background energy is also plotted for different transverse momentum thresholds, from 0 to 2 GeV/c. As one can see, while the background energy, which is proportional to \( R_c^2 \), can be reduced either reducing \( R_c \), or by applying a \( p_T \) cut (or both), the signal energy is affected by these choices only to a minor extend.

4 ALICE “First Day” Physics

After the collider closure in late 2007, the first commissioning run at \( \sqrt{s} = 900 \text{ GeV} \) is foreseen to occur at the beginning of 2008, when the ALICE setup will be characterized by the installation of the complete ITS, TPC, HMPID, Muon Arm, PMD, trigger detectors, and of partial configuration of the TOF, TRD and PHOS. Already with the first \( \sim 10^4 \) events collected during this run, some “First Day” Physics measurements will be feasible, namely the study of \( dN/d\eta \) distributions, \( p_T \) spectra, multiplicity distributions, and baryon transport analysis.

To be noted that pp runs at both the commissioning energy \( \sqrt{s} = 900 \text{ GeV} \) and at the full-energy \( \sqrt{s} = 14 \text{ TeV} \) (in late 2008) will serve not only as a baseline for the future ALICE heavy-ion program, but will also be important by themselves. As a matter of fact, the excellent expected ALICE pp performance in terms of tracking and particle identification, especially in the low \( p_T \) range, will make it complementary to the other LHC proton-proton experiments.

5 Conclusions

The ALICE experiment is characterized by a remarkable versatility in terms of the observables it will look at, and of the experimental techniques it will make use of. As the first commissioning run scheduled for the beginning of 2008 is approaching, the ALICE installation and commissioning is underway, and the experiment is evaluating its physics reach with respect to a wide variety of topics. In this report, to give some examples of the ALICE expected performance, three selected observables have been discussed,
namely heavy quarks, quarkonia and jet physics.

Acknowledgments

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HIGH DENSITY QCD PHYSICS WITH HEAVY IONS IN CMS

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The heavy ion program of the CMS experiment will examine the QCD matter under extreme conditions, through the study of global observables and specific probes.

1 Introduction

The CMS detector has a large acceptance and hermetic coverage. The various subdetectors are: a silicon tracker with pixels and strips ($|\eta| < 2.4$), electromagnetic ($|\eta| < 3$) and hadronic ($|\eta| < 5$) calorimeters, muon chambers ($|\eta| < 2.4$). The acceptance is further extended with forward detectors ($|\eta| < 6.8$). CMS detects leptons and hadrons, both charged and neutral ones. In the following, capabilities in soft, hard and forward physics are described. For a very recent extensive review see Ref. 1.

2 Soft physics

The minimum bias trigger will be based on the requirement of a symmetric number of hits in both forward calorimeters ($3 < |\eta| < 5$, see Fig. 1). For Pb-Pb collisions the centrality trigger will be provided by correlating barrel and forward energies. The charged particle multiplicity can be measured event-by-event using hits in the innermost pixel layer with about 2% accuracy and systematics below 10%.

CMS can study soft physics better than previously thought. Using a modified pixel hit triplet finding algorithm, charged particles down to very low $p_T$ can be reconstructed (Fig. 2-left). Particle identification using energy loss in silicon is possible if $p < 1-2$ GeV/c, benefitting from analogue readout. Acceptances and efficiencies are at 80-90%, the $p_T$ resolution is about 6%. At the same time low fake track rate is achieved thanks to the geometrical shape of the hit cluster: below 10% even in central Pb-Pb for $p_T > 0.4$ GeV/c. This enables the study of
identified particle spectra (down to $p_T$ of $0.1 - 0.3$ GeV/c) and yields, multiplicity distributions and correlations. Weakly decaying resonances are accessible if the found tracks are combined and selected via decay topology: strange neutral particles ($K^0_S$, Fig. 2-center, $\Lambda$, $\bar{\Lambda}$), multi-strange baryons ($\Xi^-$, $\Omega^-$). Also open charm ($D^0$, $D^{*+}$) and open beauty ($B \rightarrow J/\psi + K$) can be studied.

In Pb-Pb collisions azimuthal correlations give information on the viscosity and parton density of the produced matter. The event plane can be reconstructed using calorimetry. The estimated event plane resolution is about 0.37 rad if $b = 9$ fm. The second moment $v_2$ can be measured with about 70% accuracy. The results will improve by adding tracker information and using forward detectors, such as the zero degree calorimeter.

Figure 2: Left: Acceptance of the track reconstruction algorithm as a function of $p_T$, for tracks in the range $|\eta| < 1$. Values are given separately for pions (circles), kaons (triangles) and (anti)protons (squares). Center: Invariant mass distribution of reconstructed $K^0_S \rightarrow \pi^+ \pi^-$ in single minimum bias p-p collisions. The mass distribution of the background is indicated with a black dashed histogram. Right: Minimum bias and high level trigger $J/\psi$, $\Upsilon$, and jet trigger rates for design luminosity in central Pb-Pb collisions.
3 Hard physics

Interesting events are selected first by the level-1 trigger. It is a fast hardware trigger, decisions are made within about 3 μs after the collision. It mostly uses signals from the muon chambers and calorimeters. After that step the event rate is still high, the efficient observation of rare hard probes requires a high level trigger (HLT). The trigger uses about ten thousand CPUs working with the full event information including data from the silicon tracker. A detailed study has been done with running offline algorithms by parametrising their performance. Trigger tables are produced considering various channels and luminosity scenarios (Fig. 2-right).

Charmonium and bottomonium resonances can report on the thermodynamical state of the medium via their melting. It is an open question whether they are regenerated or suppressed at LHC energy. They can be reconstructed in the dimuon decay channel with help of precise tracking. Acceptances are at 25% (ϒ) and 1.2% (J/ψ) with 80% efficiency and 90% purity. The mass resolution is 86 MeV/ c² at the ϒ mass and 35 MeV/ c² at the J/ψ mass, in the full acceptance, and even better in the barrel (Fig. 3). This is the best resolution achieved at the LHC. With help of the HLT, 50 times more J/ψ and 10 times more ϒ will be collected.

Figure 4: Left: Expected inclusive jet $E_T$ distributions in 10 centrality bins. Right: Expected statistical reach for the nuclear modification factor for inclusive charged hadrons. For both figures, central Pb-Pb collisions at 5.5 TeV have been generated by Hydjet, with integrated luminosity of 0.5 nb⁻¹.
Finding jets on top of a high background is a challenge in Pb-Pb collisions. Jets are reconstructed using a pile-up subtraction algorithm. It consists of an iterative jet cone finder and an event-by-event background subtraction. For 100 GeV jets the directional resolutions are $\sigma_{\eta} \approx 2.8\%$, $\sigma_{\phi} \approx 3.2\%$, while the energy resolution is $\sigma_{E_T} \approx 16\%$. Thanks to the HLT, the reach of the jet $E_T$ measurement can be extended to about 0.5 TeV (Fig. 4-left). The data sets, triggered with 50, 75 and 100 GeV, are merged with a simple scaling procedure.

Parton energy loss in the hot and dense medium created in Pb-Pb collisions can be studied by measuring the nuclear modification factors $R_{AA}$ and $R_{CP}$. High $p_T$ charged particles can be tracked with about 75% algorithmic efficiency, few percent fake track rate for $p_T > 1$ GeV/c and excellent momentum resolution. Using the HLT, the $p_T$ reach of the measurement is extended from 90 to 300 GeV/c (Fig. 4-right).

4 Forward physics

The study of diffractive photoproduction of vector mesons in ultraperipheral Pb-Pb collisions can constrain the gluon density at small $x$ (Fig. 5-left). The decay channels $\rho \rightarrow \pi^+\pi^-$ and $\Upsilon \rightarrow e^+e^-$ have been studied, tagged with forward neutron detection in the zero degree calorimeter. The combined acceptance and efficiency of the method is around 20% and it gives a good mass resolution in both channels (Fig. 5-centre and right).

![Figure 5](image_url)

Figure 5: Left: The approximate $(x,Q^2)$ range covered by photoproduction in ultraperipheral Pb-Pb collisions at the LHC is indicated. Right: Invariant mass $e^+e^-$ and $\mu^+\mu^-$ distributions for photoproduced $\Upsilon$ and dilepton continuum, as expected in ultraperipheral Pb-Pb collisions at 5.5 TeV, for integrated luminosity of 0.5 nb$^{-1}$.

5 Summary

The CMS detector combines capabilities for global event characterization and for physics with specific probes. It performs equally well in soft, hard and in forward physics, often supported by high level triggering.

Acknowledgment

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References

XLI\textsuperscript{nd} RENCONTRES DE MORIOND

QCD and High Energy Hadronic Interactions

XII - Closing Session

Chairperson: J. Trân Thanh Vân
Experimental Summary Moriond QCD 2007

Gigi Rolandi

CERN,
Geneva, Switzerland

More than 90 speakers gave a presentation at this years Moriond QCD conference and more than 60 talks reported the experimental status and perspectives on Standard Model, especially QCD, search for new physics, quark spectroscopy and Heavy Ions physics. I summarize what I consider the highlights of these presentations.

1 Experimental touch

Before going into a more structured presentation I want to mention two experimental techniques that have impressed me.

1.1 CDF SVT Trigger

The CDF collaboration commissioned an impact parameter trigger based on the information provided online by the Silicon Vertex Tracker (SVT)\(^1\). The impact parameter resolution of the online reconstructed tracks is \(\sim 50 \mu m\) and allows the design of specific triggers for heavy flavor physics largely improving the efficiency in collecting b and c decays. Many results based on this trigger have been shown at this conference. Among those: the first direct measurement of CP asymmetry in the \(B_s\) system\(^2\) \(A_{CP} = 0.39 \pm 0.15\) (stat) \(\pm 0.08\) (syst) shown in Fig.1a and the angular correlation of charm pairs\(^3\) shown in Fig.1b. The latter shows that the gluon splitting plays a larger role than expected in charm pair production in ppbar collisions.

1.2 CLEO measurement of the \(D^0\) mass

The CLEO Collaboration reported\(^4\) a new measurement of the \(D^0\) mass based on \(\sim 320\) decays \(D^0 \to K_s\phi\). These events have been selected in the four millions DD pairs sample collected at the \(\Psi(3S)\). Using this rare decay channel they exploit the small Q value of the reaction (347
MeV) and the fact that the pions from the $K_s$ and the kaons from $\phi$ are in the momentum range where CLEO has optimal resolution. They obtain $M(D^0) = 1864.847 \pm 0.150\,(stat) \pm 0.095\,(syst)$ that represent an improvement of a factor 2.5 compared to the PDG value.

This is only one of the many precise measurements that CLEO-C has performed on the charm: other highlights presented at this conference are the precise measurements of exclusive branching ratios of D and D_s and the new inclusive branching ratios of D and D_s into $\eta, \eta'$ and $\phi$. Recently CLEO-C has collected 26 Millions of $\psi(2S)$ decays and the first analysis shows they can perform high precision measurements to probe light quark dynamic.

2 Parton QCD at NLO

2.1 Deep Inelastic Scattering

Generally speaking, parton QCD at NLO describes very well all existing data over many orders of magnitude. A summary of what I consider the most relevant plots is shown in Fig.2. CDF and D0 presented the inclusive jet cross section measurement showing exceptional agreement between QCD prediction and cross section in the whole explored range (p_T up to 700 GeV) when the differential cross section in y and p_T spans some 20 orders of magnitude. Here measurements in the forward region contribute to constraining the gluon Parton Distribution Functions (PDF). Also the inclusive di-jet cross section and the di-jet bbar cross section show agreement of prediction and data up to the highest di-jet mass of 400 GeV. The only distribution from the Tevatron that does not show perfect agreement between prediction and data is the jet+isolated photon cross section depending on the angle between the photon and the leading jet.

HERA collaborations have the lion share in the measurement of the DIS. The plot of the Neutral Current (NC) and Charge Current (CC) cross section in ep collisions shows perfect agreement up to the highest $Q^2$ of $3 \times 10^4$ GeV$^2$ and it is ready for the text books. It shows that NC cross section is dominated by photon exchange and that it is sensitive to the massive Z contribution at high $Q^2$; it also shows EW unification since CC and NC cross sections become similar at large $Q^2$. 
Figure 2: Various examples of comparison of DIS measurements with predictions based on QCD at NLO, see section 2.1
Tagging charm production with $D^*$, ZEUS has measured the charm production cross section as function of $Q^2$ spanning some 5 orders of magnitude and showing agreement up to the highest $Q^2$ of $\sim 500$ GeV$^2$. H1 has shown the first simultaneous measurements of the charm and beauty cross sections in photoproduction exploiting the lifetime tagging capabilities of the microvertex detector.

2.2 Gluon Content of the Proton

One of the important results of HERA is that the gluon PDFs diverge at small $x$. There is a lot of attention on gluon PDFs for their implications on QCD and also because accurate PDFs are needed for predicting cross sections at LHC. HERAII data can further constraint the gluon PDF$^8$: reduced uncertainties on high-$x$ gluons translate in reduced uncertainties in high $E_T$ jet cross section at LHC.

The low-$x$ regime at HERA has been extensively studied by H1 and ZEUS$^{10}$ also trying to understand if the DGLAP evolution works sufficiently well to extrapolate in this kinematical region and produce PDFs that can be used safely at LHC. One example of these studies is shown in Fig.3. This is the angular correlation in dijet events for various $x$. Data show disagreement with NLO at low $x$ and low $\Delta \phi$, a region where DGLAP needs checking.

Will we be able to see the saturation of the gluon PDFs at LHC$^7$? DGLAP contains a mechanism for gluon splitting, but no mechanism for non linear gluon-gluon fusions that at high density must balance the branching$^{11}$. The LHC forward rapidity region allows to reach very small $x : x \sim 10^{-6}$ for $y \sim 6$ and $Q \sim 10$ GeV and possibly study the gluon saturation. The CASTOR detector at CMS covers rapidity up to 6.5.

3 Heavy Ions Physics

RICH data have shown that the particle multiplicity in Heavy Ions (HI) collisions at $\sqrt{s} = 200$ GeV is at the low edge of the predictions before RICH-era. This also implies that tracking and other measurements in HI collisions at LHC will not be as difficult as anticipated. RICH data show$^{12}$ that particle production in HI collisions follows a quite simple pattern since the distributions of many geometrical variables like pseudorapidity and elliptic flow scale with energy in the whole energy range explored by RICH.
The STAR collaboration reported evidence for jet quenching in HI collisions: the presence of back to back particle correlation in d-Au collisions and its absence in Au-Au collisions show that in Au-Au jets loose energy when they cross the dense colored matter. At this conference STAR reported that when studying the hadrons correlations in HI collisions one has to distinguish different regions in the $\Delta \eta - \Delta \phi$ plane (see Fig.4): there is a near side jet peak, an away side and also a near side ridge that includes particle produced at small $|\Delta \phi|$ and all $\Delta \eta$ in the wake of the longitudinal flow. The ridge is correlated to jet production and persists to the highest $p_T$. Subtracting the effects of the ridge recovers centrality independent jet yields.

4 Electroweak Physics

4.1 Precise measurement of the $W$ mass

CDF reported a new and precise measurement of the $W$ mass and width:

$$M_W = 80417 \pm 48(\text{stat} + \text{syst}) \text{ MeV}$$

In order to achieve the remarkable precision of 0.06% CDF studies the transverse mass distribution in semileptonic $W$ decays and utilizes well understood data samples to calibrated the detector to high precisions. The momentum scale and resolution are given by the peaks of the $\Upsilon(1S)$ and $Z^0$ in the di-muon channel. This scale is ported to the calorimeter studying the $E/p$ ratio with high momentum electrons and checked using $Z^0$ decays into two electrons. The momentum of the neutrino is obtained from missing energy and the systematic effects are studied using events where $Z$ decays in two muons. The two measurements in the di-electrons and di-muons channels are very compatible and they are averaged to obtain the final result. The momentum scale gives the main systematic error. This is the most precise single measurement of the $W$ mass and improves the precision of the world average reducing its uncertainty by 15%.

4.2 Precise measurement of the top mass

CDF and D0 reported new measurements of the $t\bar{t}$ cross section and of the top mass. Though many different analyses exploiting all decay channels contribute to the final value of the mass (see Fig.5), the most precise value is obtained from events where one top decays semileptonically and
the other in jets (lepton + jets channel) that correspond to \( \sim 30\% \) of the \( \bar{t}t \) events. The analyses used to determine the top mass are quite sophisticated, often requiring huge computations. The jet energy scale is fitted together with the top mass, using the known mass of the \( W \). The main systematic error is the modeling of the initial and final state radiation. All methods extract the mass comparing data with simulation and the measured mass is the mass as defined in PYTHIA. In the near future we should invest some time to understand in detail if this definition has any implication for the use of the top mass in the electroweak fit. This was discussed quite lively at the conference, but no final conclusion was reached. The present value of the top mass is

\[
m_{\text{top}} = 170.9 \pm 1.1 \text{(stat)} \pm 1.5 \text{(syst)} \text{ GeV}
\]

with a remarkable precision of 1.1%.

4.3 Electroweak Fit

Figure 6a shows the "blue-band plot", the \( \Delta \chi^2 \) curve derived from high \( Q^2 \) precision electroweak measurements, performed at LEP and by SLD, CDF and D0, as a function of the Higgs-boson mass, assuming the Standard Model to be the correct theory of nature. It includes now the precise measurements presented at this conference.

The minimum of the curve is today at 76 GeV, with an experimental uncertainty of \( +33 \) and \(-24 \) GeV (at 68% confidence level derived from \( \Delta \chi^2 = 1 \) for the black line, thus not taking the theoretical uncertainty shown as the blue band into account). The precision electroweak measurements tell us that the mass of the Standard-Model Higgs boson is lower than about 144 GeV (one-sided 95% confidence level upper limit derived from \( \Delta \chi^2 = 2.7 \) for the blue band, thus including both the experimental and the theoretical uncertainty). This limit increases to 182 GeV when including the LEP-2 direct search limit of 114 GeV shown in yellow.

The preferred value of the Higgs Boson mass from the Electroweak fit has evolved with time (see Fig. 6b), following the improved precision of the electroweak measurements. We notice that now - for the first time - the upper side error band of the Higgs mass is below the direct search limit.

There are different ways to look at this plot. One way is to say that the Standard Model is correct and that the probable range for the Higgs mass is now quite narrow: \( 114 \text{ GeV} < M_H < \)
Figure 6: a) $\Delta \chi^2$ of the Electroweak fit as function of the Higgs mass. b) Evolution with time of the fitted value of the Higgs mass from the Electroweak fit: squares 95% confidence limit, dots Higgs Mass.

Figure 7: The measurements of $M_W$, $m_{top}$ and $\sin^2 \theta_{ew}$ are compared with the prediction of Standard Model (SM) and Minimal SuperSymmetry Model (MSSM).
144 GeV. Another way is to recall that the plot describes the Higgs Mass only in the assumption that the Standard Model is correct: if we lift this hypothesis there are other contributions to the radiative corrections ($\Delta \rho$) that have to be taken into account. One example is given in the plots shown in Fig. 7 where the measurements of $M_W$, $m_\text{top}$ and $\sin^2\theta_\text{eff}$ are compared with the prediction of Standard Model (SM) and Minimal SuperSymmetry Model (MSSM). We see that the SM barely fits the data at 68% confidence level while there are MSSM regions that are well compatible with all measurements.

5 Higgs Boson search, single top and other known rare phenomena

The electroweak data suggest to search the SM Higgs in the low mass range ($M_H < 160$ GeV) that is accessible at the TEVATRON, albeit with small cross sections ($< 1$ fb). CDF and D0 have presented new results. Since the cross sections are small the search is done in the Higgs decay channels with large branching ratio: $b\bar{b}$ for $M_H \leq 135$ GeV and WW for $M_H \geq 135$ GeV.

This is different from what is foreseen at LHC where the Higgs cross section and the luminosity are larger. Thanks to the larger production rate the Higgs is searched for in more clean channels with small branching ratio ($\gamma\gamma$ or $\tau\tau$ in Higgs produced via Vector Boson Fusion) that give better sensitivity.

TEVATRON searches exploit three main channels for "low" mass Higgs: $ZH \to \ell\ell bb$, $ZH \to \nu\nu bb$ and $WH \to t\bar{t}bb$. And one channel for "high" mass Higgs: $gg \to H \to WW \to t\bar{t}t\bar{t}$. Finding the Higgs is possible but very challenging: with present analyses in 1 fb$^{-1}$ the number of selected events in each channel after basic acceptance cuts ranges typically from 1 to 4 events while the background ranges from 100 to 500 events. CDF and D0 have started a campaign to improve signal acceptance and reduce the backgrounds whenever possible. The present sensitivities with 1 fb$^{-1}$ range between 5 times and 10 times the standard model cross section, depending on the Higgs mass. The observed limits (see Fig. 8) are similar to the expected ones for Higgs masses around 160 GeV and somewhat higher for lower masses. The main question is how to extrapolate to the full luminosity expected to be collected by the Tevatron experiments (8 fb$^{-1}$): the region near 160 GeV seems accessible, while low masses
require a very important improvement of the analyses.

The analyses done to observe some rare expected phenomena in the present data sets are an important benchmark for the Higgs searches. The D0 collaboration presented evidence for the single top production at a level of 3.4σ. They measure a cross section that is higher, but compatible, with the Standard Model. This is not unusual for a first observation where one profits of an upward statistical fluctuation. For the same reason it is not unusual that single top production has not yet been seen by CDF with similar amount of statistics. CDF presented evidence of ZZ production in the channel $ZZ \rightarrow \ell\ell(\ell\ell + \nu\nu)$.

In these analyses the number of the signal events after basic acceptance cuts is 36 with a background of 600 for the single top and of 9 with a background of 23 for the ZZ. The final extraction of the signal is done exploiting efficient statistical techniques like Boosted Decision Trees or Matrix Element methods that use all the characteristics of the selected events. These techniques will be used also for the extraction of the Higgs signal. One example is the new CDF analyses for the Higgs decay in WW.

6 Heavy Flavor Physics and Unitarity Triangle

The experimental tests of the unitarity triangle have been reviewed excellently by Stocchi. The world average of $\sin(2\beta_{\text{eff}})$ in the golden channel $b \rightarrow ccs$ is now $\sin(2\beta_{\text{eff}}) = 0.68 \pm 0.03$.

The two sides of the unitarity have lengths $|V_{ub}|/|V_{cb}|$ and $|V_{ts}|/|V_{td}|$. The new precise measurement of $\Delta m_s$ at the Tevatron is combined with the known value of $\Delta m_d$ to obtain an indirect determination of $|V_{ts}|/|V_{td}|$ of $0.208 \pm 0.07$. The error on $V_{ts}^2/V_{td}^2$ (6%) is completely determined by uncertainties in the theory while the experimental errors on $\Delta m_s$ and $\Delta m_d$ are 0.7% and 1% respectively. There is now progress on the determination $|V_{ub}|$ based on the measurements of the form factors in exclusive channels and on studies of the stability in the extrapolation to the full phase space of the inclusive measurements. The sides of the triangle can be used to obtain an indirect determination of $\sin(2\beta_{\text{eff}})$. This is actually done making a global fit to all available information but from the direct measurement of $\sin(2\beta_{\text{eff}})$. The result of the fit is $\sin(2\beta_{\text{eff}}) = 0.76 \pm 0.04$ that is somewhat higher (1.6σ) than the direct measurement.
The golden channel $b \to ccs$ provides a precise and direct measurement of $\sin(2\beta^{eff})$. Many other channels $b \to qqs$ can be used for less direct determinations of $\sin(2\beta^{eff})$ (see Fig.9) that are affected by the uncertainties of diagrams involving penguins. The error bars presented in Fig.9 include the theoretical errors. The naive average of all these other determination is $\sin(2\beta^{eff}) = 0.53 \pm 0.05$ that is somewhat lower (2.6$\sigma$) than the direct measurement. This discrepancy should not be taken too seriously because the naive average is not the correct way to combine all these different measurement. It would be interesting to see if there are correlations among the interpretations of the different indirect determinations and if these correlations can be computed and expressed in terms of few variables. I look forward to an analysis of the different reactions similar to what has been done in the electroweak sector by Peskin et al.\cite{27} or Altarelli et al.\cite{28} who have shown that the different experimental observables can be interpreted with a small number of model independent parameters.

7 D$^0$ $\bar{D}^0$ mixing

BaBar presented the first evidence of D$^0$ $\bar{D}^0$ mixing\cite{29}. Mixing has been seen studying the time dependence of the ratio $R_{WS}$ between the wrong sign and the right sign decay yields of the $D^0 \to K\pi$ (see figure 10). The sign of the charged kaon is compared to the sign of the slow $\pi$ in $D^* \to D^0\pi$ decays and the time is measured from the decay length of the $D^0$. The time dependence contains three terms: a constant term for the Double Cabibbo Suppressed (DCS) decays, a quadratic term for the oscillations and a linear term for the interference between the two amplitudes.

$$R_{WS} = R_{DCS} + \sqrt{R_{DCS}} y'\left(\frac{t}{\tau}\right) + \frac{x'^2 + y'^2}{2} \left(\frac{t}{\tau}\right)^2$$

where $\tau$ is the $D^0$ lifetime.

The two terms $x'$ and $y'$ are linked to $x = \Delta m/\Gamma$ and $y = \Delta \Gamma/(2\Gamma)$ by a rotation

$$x' = x \cos \delta + y \sin \delta \quad y' = -x \sin \delta + y \cos \delta$$

involving a strong phase $\delta$.

By fitting the time dependence BaBar finds a 3.9$\sigma$ evidence of the oscillations with $y' = (9.7 \pm 4.4\pm 3.1) \times 10^{-3}$ and $x'^2 = (-0.22 \pm 0.30\pm 0.20) \times 10^{-3}$ with important correlation between the two variables.

Belle presented\cite{30} a time dependent analysis of the Dalitz plot in the decay $D^0 \to K^0_S\pi^+\pi^-$ and measures $x = (0.80 \pm 0.29\pm 0.17) \times 10^{-2}$ and $y = (0.33 \pm 0.24\pm 0.15) \times 10^{-2}$ also with large correlation between the two variables because this time dependence essentially measures $x^2 + y^2$. 

Figure 10: Time dependence of the ratio $R_{WS}$ between wrong and right sign $D^0$ decays.
Belle measured also the lifetime difference between $D^0 \to K^- \pi^+$ and $D^0 \to K^+ K^-, \pi^+ \pi^-$. In the assumption that the CP violation effects are negligible this provides an independent measurement of $y = (1.31 \pm 0.32 \pm 0.25) \times 10^{-2}$.

This information can be analyzed in the standard model giving a consistent picture that allows also a determination of the strong phase $\delta = -38^\circ \pm 46^\circ$.

8 Conclusions

Very large samples of data are available allowing precise measurements and studies of rare phenomena. I have been impressed by the quality and the complexity of the analyses that have been presented. The steady progress in the field is marked by discoveries of new expected phenomena that have been recently seen: D̅D̅ mixing, single top, ZZ.

QCD is - generally speaking - in good shape and describes well the hadron production at high and low $Q^2$. This accurate description requires a monumental effort both on theory and experimentation and is needed for the searches of new physics at hadron colliders and at lower energy.

Everybody is waiting for the start of LHC. It is close. At Moriond XLIV we will hopefully see the first presentations with data, but it will take time to understand the detectors. In the mean time watch the Tevatron, the Triangle and the rest.... The devil is in the details of experiments and theory.

I would like to give my warmest thank to Tran for the 42 Moriond that he invented for us, to the organizers for a well organized conference and to all the speakers for their very good talks.

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MORIOND QCD 2007 – THEORY SUMMARY

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Developments reported at the 2007 Moriond Workshop on QCD and Hadronic Interactions are reviewed and placed in a theoretical context.

1 Introduction

QCD was invented in 1973. (There were some earlier hints.) We are still concerned with it as neither perturbative nor currently available non-perturbative (e.g., lattice) methods apply to many interesting phenomena. These include hadron structure, spectroscopy, jet and quarkonium fragmentation, heavy ion physics, and effects of thresholds. The understanding of hadronic behavior is crucial in separating underlying short-distance physics (whether electroweak or new) from strong-interaction effects. The properties of hadrons containing heavy quarks provide an exceptional window into QCD tests. Finally, QCD may not be the only instance of important non-perturbative effects; familiarity with it may help us to prepare for surprises at the Large Hadron Collider (LHC). In this review we shall discuss a number of developments reported at Moriond QCD 2007 in the context of these ideas. A companion review 1 deals directly with the experimental results. I apologize for not covering some theoretical topics whose relation to experimental results presented at this conference is not yet clear to me, and for omitting some nice experimental results for which I have no comments.

2 Heavy flavor issues: the current CKM matrix

The Kobayashi-Maskawa matrix theory of CP violation, and its parametrization of charge-changing weak transitions, as shown in Fig. 1, passes all experimental tests so far. The major uncertainties in the parameters of the CKM matrix are now dominated by theory. Briefly, we have $V_{ud} ≈ V_{us} ≈ 0.974$, $V_{us} ≈ -V_{cd} ≈ 0.226$, $V_{cb} ≈ -V_{ts} ≈ 0.041$, $V_{td} ≈ 0.008 e^{-i 21^\circ}$, ...
$V_{ub} \approx 0.004 e^{-i \theta_{ub}}$ (sources of phase information will be explained below), and – on the basis of single-top production observed by the D0 collaboration $^2$ $-0.68 < |V_{td}| < 1$ at 95% c.l.

3 Meson decay constants and implications

The ability of theory to anticipate important hadronic properties is illustrated by recent results on meson decay constants. Moreover, it has been possible in some cases to replace calculated quantities with better-determined experimental ones, reducing errors on fundamental parameters such as CKM matrix elements.

In 2005 the CLEO Collaboration $^3$ reported the measurement $f_{D^+} = (222.6 \pm 16.7^{+2.8}_{-3.4})$ MeV, to be compared with one lattice QCD prediction $^4$ of $201 \pm 3 \pm 17$ MeV. More recently CLEO has measured $f_{D_s} = (274 \pm 13 \pm 7)$ MeV. $^5$ One can obtain a slightly more precise value by including preliminary data on $D_s \rightarrow \tau \nu$ where $\tau \rightarrow e \nu \bar{\nu}$. $^6$ The BaBar Collaboration reports $f_{D_s} = (283 \pm 17 \pm 14)$ MeV. $^7$

One lattice prediction $^4$ is $f_{D_s} = 249 \pm 3 \pm 16$ MeV, leading to a predicted ratio $f_{D_s}/f_D = 1.24 \pm 0.01 \pm 0.07$. This is to be compared with the CLEO ratio $1.23 \pm 0.11 \pm 0.04$. $^5$ One expects $f_{B_s}/f_B \approx f_{D_s}/f_D$ so better measurements of $f_{D_s}$ and $f_D$ by CLEO will help validate lattice calculations and provide input for interpreting $B_s$ mixing. A desirable error on $f_{B_s}/f_B \approx f_{D_s}/f_D$ is $\leq 5\%$ for a useful determination of the CKM element ratio $|V_{td}/V_{ts}|$. This will require errors $\leq 10$ MeV on $f_{D_s}$ and $f_D$. (Independent information on $|V_{td}/V_{ts}|$ has come from a precise measurement of $B_s \rightarrow \tau \nu$ mixing. $^8$) A scaling argument from the quark model $^9$ implies $f_{D_s}/f_D \approx f_{B_s}/f_B \approx \sqrt{M_s/M_d} \approx 1.25$, with constituent masses $M_s \approx 485$ MeV, $M_d \approx 310$ MeV.

4 $B_s$ physics

Comparing box diagrams for $b \bar{s} \rightarrow s \bar{b}$ and $bd \rightarrow db$ (dominated by intermediate top quarks), one sees that $B_s \rightarrow \bar{B_s}$ mixing is stronger than $B \rightarrow \bar{B}$ mixing because $|V_{ts}/V_{td}| \approx 5$. Now, CKM unitarity implies $|V_{ts}| \approx |V_{td}| \approx 0.041$ is well measured, so $B_s \rightarrow \bar{B_s}$ mixing really probes the matrix element between $B_s$ and $\bar{B_s}$. This quantity involves $f_{B_s} B_{B_s}$, whose ratio with respect to that for non-strange $B$’s is known from lattice QCD: $^{10}$

$\xi \equiv f_{B_s} \sqrt{B_{B_s}}/(f_B \sqrt{B_B}) = 1.21 \pm 0.02$. $^{10}$

The $B^0 - \bar{B}^0$ mixing amplitude is well-measured: $\Delta m_d = (0.507 \pm 0.004)$ ps$^{-1}$. Consequently,
Figure 2: Examples of decay topologies for $B^0 \rightarrow \pi^+\pi^-$. (a) Tree, (b) penguin.

measurement of $B_s$ mixing implies a value of $|V_{cd}/V_{ts}|$. The recent CDF measurement at Fermilab
$\Delta m_s = (17.77 \pm 0.10 \pm 0.07) \text{ ps}^{-1}$ gives $|V_{cd}/V_{ts}| = 0.206 \pm 0.008$ and hence $1 - \rho - \eta \equiv |V_{ub}^* V_{td}/(V_{cb}^* V_{cd})| = 0.91 \pm 0.04$. This implies that $\gamma \equiv \text{Arg}(V_{ub}^* V_{td}/(V_{cb}^* V_{cd})) \simeq (66 \pm 6)^\circ$, a great improvement over previous determinations.

The first evidence for $B_s$ mixing was presented by the D0 collaboration. This collaboration
has now presented evidence for a decay rate difference between the $B_s$ mass eigenstates, with
the eigenstate which is approximately CP-even decaying somewhat more rapidly: $\Delta \Gamma_s = 0.13 \pm 0.09 \text{ ps}^{-1}$. This agrees with the expected value $\Delta \Gamma_s \approx (1/200) \Delta m_s \approx 0.09 \text{ ps}^{-1}$. (The
values of $\Delta \Gamma_s$ and $\Delta m_s$ are expected to track one another.) Within large errors, D0 sees no
evidence for CP violation in $B_s \rightarrow J/\psi \phi$. One expects in the Standard Model $\phi_s = 0.036 \pm 0.003$, a
correct which may be accessible to LHCb.\textsuperscript{14,15}

5 Systematics of $B$ decays

5.1 General considerations

Reviews of $B$ decays were given at this Conference by Lim\textsuperscript{16} (experiment) and Lu\textsuperscript{17} (theory).
It is useful to visualize $B$ decay amplitudes in terms of flavor diagrams (see, e.g., Fig. 2). Flavor SU(3) permits one to relate decay asymmetries in one channel to those in another. For
example, one can show\textsuperscript{19,20}

$$\Gamma(\bar{B}^0 \rightarrow \pi^+\pi^-) - \Gamma(B^0 \rightarrow \pi^+\pi^-) = -\frac{1}{2} \Gamma(B^0 \rightarrow K^-\pi^+) - \Gamma(B^0 \rightarrow K^+\pi^-) \right) .$$

(1)

Using dominance of $B \rightarrow K\pi$ transitions by the isospin-preserving ($\Delta I = 0$) penguin $\bar{b} \rightarrow s$ transition, and a well-established hierarchy of other amplitudes, one can obtain sum rules for rates\textsuperscript{21} and asymmetries\textsuperscript{22,23} in these decays. Defining the CP-averaged ratios

$$R = \frac{\bar{\Gamma}(B^0 \rightarrow K^+\pi^-)}{\bar{\Gamma}(B^+ \rightarrow K^+\pi^-)} , \quad R_c = \frac{2\bar{\Gamma}(B^0 \rightarrow K^+\pi^0)}{\bar{\Gamma}(B^+ \rightarrow K^0\pi^+)} , \quad R_n = \frac{\bar{\Gamma}(B^0 \rightarrow K^+\pi^-)}{2\bar{\Gamma}(B^0 \rightarrow K^0\pi^0)}$$

(2)

where $\bar{\Gamma}(B \rightarrow f) \equiv [\Gamma(B \rightarrow f) + \Gamma(\bar{B} \rightarrow \bar{f})]/2$, one such sum rule is $R_c = R_n$. Experimentally\textsuperscript{24}

$$R = 0.90 \pm 0.05 , \quad R_c = 1.11 \pm 0.07 , \quad R_n = 0.97 \pm 0.07 ,$$

(3)

so the sum rule is satisfied. It is expected to hold also to first order in isospin breaking.\textsuperscript{25}

A recent result is relevant to the systematics of $B \rightarrow PV$ decays, where $P$ and $V$ are light pseudoscalar and vector mesons. The pure penguin process $B^+ \rightarrow K^0\rho^+$ has been seen by BaBar\textsuperscript{26} with a branching ratio $B(B^+ \rightarrow K^0\rho^+) = (8.0^{+1.4}_{-1.3} \pm 0.5) \times 10^{-6}$. This is comparable
to the pure-penguin process $B^+ \to K^{*0} \pi^+$ with $B = (10.7 \pm 0.8) \times 10^{-6}$. In the former process, the spectator quark ends up in a vector meson ($^3p^-$), while in the latter the spectator ends up in a pseudoscalar ($^3p^-$). This confirms an early expectation by Lipkin that the amplitudes for the two processes were related by $p^- \simeq -p^-$.  

5.2 $B_s$ decays

One way to learn the width difference $\Delta \Gamma$ of $B_s$ mass eigenstates is to compare the decay lifetimes in different polarization states of the final vector mesons in $B_s \to J/\psi \phi$. These are conveniently expressed in a Cartesion basis. There are three such states. Two are CP-even. In one of these, the vector mesons' linear polarizations are perpendicular to the decay axis and parallel to one another ("||"). In the other CP-even state, both vector mesons are longitudinally polarized ("||"). In the CP-odd state, the vector mesons' linear polarizations are perpendicular to the decay axis and also to one another ("\perp"). Separating out the CP-even and CP-odd lifetimes would be much easier using $\parallel$ and $\perp$ states, thereby avoiding bias due to imperfect modeling of polar angle dependence.

The branching ratio $B(B_s \to K^+ K^-) = (24.4 \pm 1.4 \pm 4.6) \times 10^{-6}$ reported by CDF at this Conference is due mainly to the $|\Delta S| = 1$ penguin. For comparison, $B(B^+ \to K^{0} \pi^+) = (23.1 \pm 1.0) \times 10^{-6}$. The large error on the former means that one can't see the effects of non-penguin amplitudes through interference with the dominant penguin.

$B_s$ decays help validate flavor-SU(3) techniques used in extracting CKM phases. For example, under the U-spin transformation $d \leftrightarrow s$, the decay $B_s \to K^- \pi^+$ is related to $B^0 \to K^+ \pi^-$. It has a branching ratio of $(5.0 \pm 0.75 \pm 1.0) \times 10^{-6}$; it differs from the process $B^0 \to \pi^+ \pi^-$ with $B = (5.16 \pm 0.22) \times 10^{-6}$ only by having a different spectator quark.

5.3 Baryonic $B$ decays

Results presented at this conference shed light on the mechanisms of $B$ decays to baryonic final states. Low-mass baryon-antibaryon enhancements seen in these decays favor a fragmentation picture over resonant substructure, based in part on information from angular correlations between decay products. The production of several heavy quarks, as in $b \to c \bar{c}$, helps produce baryons like $c s q$ where $q = (u, d)$ gives $\Xi_c$ and $q = s$ gives $\Omega_c$. The large available phase space and high quark multiplicity in $B$ decays may permit the production of exotic final states.

5.4 Sum rules for $CP$ asymmetries in $B \to K \pi$

Using the dominance of the $\Delta I = 0 \ b \to \bar{s}$ penguin amplitude, M. Gronau has shown that

$$A_{CP}(K^+ \pi^-) + A_{CP}(K^{0} \pi^+) = A_{CP}(K^+ \pi^0) + A_{CP}(K^{0} \pi^0).$$

(4)

Non-penguin amplitudes should be small in $B^+ \to K^{0} \pi^+$, so $A_{CP}(K^{0} \pi^+) \simeq 0$ and

$$A_{CP}(K^+ \pi^-) = A_{CP}(K^+ \pi^0) + A_{CP}(K^{0} \pi^0).$$

(5)

[Strictly speaking, a more accurate version of these sum rules applies to CP-violating rate differences $\Delta(I) = \Gamma(B \to \bar{f}) - \Gamma(B \to f).$] The observed CP asymmetries are $A_{CP}(K^+ \pi^-) = -0.097 \pm 0.012$, $A_{CP}(K^{0} \pi^+) = 0.009 \pm 0.025$, $A_{CP}(K^+ \pi^0) = 0.047 \pm 0.026$, and $A_{CP}(K^{0} \pi^0) = -0.12 \pm 0.11$. The last is the most poorly known and may instead be predicted using the sum rules. With corrections for $\tau(B^+)/\tau(B^0) = 1.076 \pm 0.008$ and branching ratios, the first and second of these sum rules predict $A_{CP}(K^{0} \pi^0) = (-0.140 \pm 0.043, -0.150 \pm 0.035$. The experimental value of $A_{CP}(K^{0} \pi^0)$ carries too large an error at present to provide a test.
A vanishing $A_{CP}(K^0\pi^0)$ would imply $A_{CP}(K^+\pi^-) = A_{CP}(K^+\pi^0)$, which is not so. $A_{CP}(K^+\pi^0)$ and $A_{CP}(K^0\pi^0)$ involve color-suppressed tree and electroweak penguin (EW) amplitudes. The latter occur in a calculable ratio $\delta_{EW} = 0.60 \pm 0.05$ with respect to known amplitudes.

One may ask how the CP asymmetry in $B^0 \rightarrow K^+\pi^-$ can be non-zero, thereby signaling the presence of non-penguin amplitudes, while neither the CP asymmetry nor the rate ratio $R_c$ shows evidence of such amplitudes in $B^+ \rightarrow K^+\pi^0$. Let $r_c \sim 0.2$ denote the ratio of tree to penguin amplitudes in $B^+ \rightarrow K^+\pi^0$. One may write the sum rule

$$\left( \frac{R_c - 1}{\cos \gamma - \delta_{EW}} \right)^2 + \left( \frac{A_{CP}(B^+ \rightarrow K^+\pi^0)}{\sin \gamma} \right)^2 = (2r_c)^2 + \mathcal{O}(r_c^3),$$

which is essentially based on the identity $\cos^2 \delta + \sin^2 \delta = 1$, where $\delta$ is a strong phase. The key to this sum rule's validity is that $\cos \gamma \simeq \delta_{EW}$, thereby allowing it to be satisfied for $R_c \sim 1$ and small $A_{CP}(K^+\pi^0)$.

5.5 Ways to measure $\sin2\beta$

The BaBar Collaboration has updated its value based on $b \rightarrow cs\bar{s}$ decays, $\sin2\beta = 0.714 \pm 0.032 \pm 0.018$. When combined with the latest Belle value, $\sin2\beta = 0.642 \pm 0.031 \pm 0.017$ and earlier data this gives a world average $\sin2\beta = 0.678 \pm 0.025$, serving as a reference for all other determinations of $\beta$.

Recently BaBar studied in the decay $B^0 \rightarrow D^{(*)0}_s h^0_s$, extracting coefficients $S$ and $C$ of time-dependent decay rate modulations proportional to $\sin \Delta m t$ and $\cos \Delta m t$. The result $\sin 2\beta_{\text{eff}} = -S = 0.56 \pm 0.23 \pm 0.05$ is compatible with the reference value. The value $C = -0.23 \pm 0.16 \pm 0.04$ is compatible with no direct CP violation, as expected in the Standard Model, but carries a large experimental error.

A large number of processes are dominated by $b \rightarrow s$ penguin amplitudes. When averaged, these give $\sin 2\beta_{\text{eff}} = 0.53 \pm 0.05$, a value 2.6$\sigma$ below the reference value. It is not clear that it makes sense to average all these processes as some involve $b \rightarrow s\bar{s}s$, others $b \rightarrow s\bar{d}d$ and/or $b \rightarrow s\bar{t}u$, and some involve mixtures. Moreover, QCD corrections can differ for different final states. The experimental values have shifted a good deal from year to year, providing theorists with a moving target which they have been quite adept at following. At present the number on which I am keeping an eye is that from $B^0 \rightarrow \pi^0 K_S$, which both BaBar and Belle agree lies below the reference value, with an average $\sin 2\beta_{\text{eff}} = 0.33 \pm 0.21$. (Note the large experimental error.) The value of the $\cos \Delta m t$ coefficient $C_{K\pi\pi}$ is $0.12 \pm 0.11$ also interesting. This is just $-A_{CP}(K^0\pi^0)$. As noted earlier, sum rules predict a central value of 0.14 to 0.15 for $C_{K\pi\pi}$.

Many estimates have been performed of deviations of $\sin 2\beta_{\text{eff}}$ from the reference value in the Standard Model. Typical explicit calculations give a deviation of 0.05 or less, usually predicting $\sin 2\beta_{\text{eff}}$ larger than 0.678 whereas most experiments find lower values. Flavor-SU(3) estimates allow differences of at most 0.1.

5.6 CP violation in $B \rightarrow \pi\pi$

An example of the systematic error associated with uncertainty in hadron physics is provided by a detailed examination of time-dependent CP asymmetries in $B^0 \rightarrow \pi^+\pi^-$. This is relevant to remarks made by Liu at this Conference concerning limitations in our ability to learn the weak phases $\alpha$ and $\gamma$. I report on work with M. Gronau, updating a previous analysis.

The time-dependent asymmetry parameters $(S_{\pi\pi}, C_{\pi\pi})$ have been measured by BaBar $(-0.60 \pm 0.11, -0.21 \pm 0.09)$ and Belle $(-0.61 \pm 0.11, -0.55 \pm 0.09)$, leading to an average $(0.605 \pm 0.078, -0.376 \pm 0.066)$. These average values are plotted in Fig. 3 along with predictions for values of the weak phase $\alpha$ and strong phase $\delta = \delta^P - \delta^T$. An SU(3)-breaking factor
\( f_K/f_\pi = 1.22 \) has been taken for the ratio of \(|\Delta S| = 1\) to \(\Delta S = 0\) tree amplitudes, but no SU(3) breaking has been assumed for the corresponding ratio of penguin amplitudes. The error ellipses represented by the plotted points encompass the ranges \(81^\circ \leq \alpha \leq 91^\circ\) (implying \(68^\circ \leq \gamma \leq 78^\circ\)) and \(-40^\circ \leq \delta \leq -26^\circ\). As in Ref.\(^{44}\), we get a very small range of \(\gamma\) [here \((73 \pm 4)^\circ]\), but additional systematic errors are important. In the upper figure, the penguin "pollution" has been estimated using \(B^+ \to K^0\pi^+\), entailing the neglect of a small "annihilation" amplitude, while in the lower figure it has been estimated using \(B^0 \to K^+\pi^-\), in which the effect of a small tree amplitude must be included. The two methods give weak phases within a degree or two of one another.

Now we examine the effect of SU(3) breaking in the ratio of penguin amplitudes. Call the \(\Delta S = 0\) penguin \(P\), and the \(|\Delta S| = 1\) penguin \(P'\), and define \(\xi_P \equiv |P^\prime/P|V_{ub}^*V_{ub}/V_{us}^*V_{us}\). The above exercise was for \(\xi = 1\). Now we vary \(\xi_P\).

One could assume \(\xi_P = f_K/f_\pi = 1.22\) as for the tree amplitude ratio.\(^ {44}\) Alternatively, one could determine it from \(\Delta(K^+\pi^-) = -\xi_P\Delta(\pi^+\pi^-)\), where \(\Delta(f) \equiv \Gamma(B^+ \to f) - \Gamma(B \to f)\). In this case with the world average \(A_{CP}(K^+\pi^-) = -0.097 \pm 0.012\) one finds \(\xi_P = 0.79 \pm 0.18\).

The change from \(\xi = 1\) to \(\xi_P = 1.22\) shifts \(\alpha\) up \((\gamma\) down) by \(\sim 8^\circ\), \(|\delta|\) up by \(\sim 10^\circ\), while the change to \(\xi_P = 0.79\) shifts \(\alpha\) down \((\gamma\) up) by \(\sim 10^\circ\), \(|\delta|\) down by \(\sim 8^\circ\). The systematic (theory) errors are larger than the statistical ones. As stressed by Lü,\(^ {17}\) one needs to gain control of SU(3) breaking. In order to provide information beyond that obtained from flavor SU(3), schemes such as PQCD\(^ {17}\) and SCET\(^ {45}\) need to predict \(\delta\) to better than \(10^\circ\).

Discussion at this Conference concerned the relative merits of frequentist\(^ {46}\) and Bayesian\(^ {47}\) analysis, referring to a recent controversy over what can be learned from \(B \to \pi\pi\).\(^ {48}\) The intelligent choice of priors can have merits, e.g., when searching for a point on the surface of a sphere (taking a uniform prior in the cosine of the polar angle \(\theta\), not \(\theta\) itself) or when searching for a lost skier at La Thuile (beginning by looking near the lifts).
6 \ D \ \text{mixing}

In the Standard Model, mixing due to shared intermediate states reached by $|\Delta C| = 1$ transitions dominates $D^0$-$\bar{D}^0$ mixing. In the flavor-SU(3) limit these contributions (e.g., $\pi\pi$, $K\bar{K}$, $K\pi$, and $\bar{K}\pi$) cancel one another. How precise is the cancellation?

Define $D_1$ and $D_2$ to be the mass eigenstates (respectively CP-even and -odd in the absence of CP violation), $\Delta M \equiv M_1 - M_2$, $\Delta \Gamma \equiv \Gamma_1 - \Gamma_2$, $x \equiv \Delta M/\Gamma$, and $y \equiv \Delta \Gamma/\Gamma$, where $\Gamma \equiv (\Gamma_1 + \Gamma_2)/2$. Estimates of $y$ range up to $O(1\%)$, with $|x| \leq |y|$ typically.

The time dependence of "wrong-sign" $D^0(t = 0)$ decays (e.g., to $K^+\pi^-$) involves the combinations $\epsilon' = \epsilon \cos \delta_{K\pi} + y \sin \delta_{K\pi}$, $y' = -\epsilon \sin \delta_{K\pi} + y \cos \delta_{K\pi}$, where the strong phases $\delta_{K\pi}$ has been measured by the CLEO Collaboration: $\delta_{K\pi} = 1.09 \pm 0.66$. In the SU(3) limit, $\delta_{K\pi} = 0$. This method has been used by the BaBar Collaboration to obtain the non-zero mixing parameter $y' = (9.7 \pm 4.4 \pm 3.1) \times 10^{-3}$.

The Belle Collaboration has obtained evidence for mixing in a different way, by comparing lifetimes in CP- and flavor-eigenstates and thereby measuring a parameter $y_{CP} = (1.13 \pm 0.32 \pm 0.25)\%$. In the limit of CP conservation (a likely approximation for $D$ mesons), $y_{CP} = y$. A time-dependent Dalitz plot analysis of $D^0 \to K\pi\pi^0$ by Belle obtains $x = (0.80 \pm 0.29^{+0.09+0.15}_{-0.04-0.14})\%$, $y = (0.33 \pm 0.24^{+0.07+0.08}_{-0.12-0.06})\%$.

These results were synthesized in several theoretical analyses. The consensus is that while $y$ is near the upper limit of what was anticipated in the Standard Model, there is no evidence for new physics. Observation of CP violation in $D$ decays, on the other hand, would be good evidence for such physics, and will continue to be the object of searches.

7 \ Low-energy hadron physics

Information on light-quark interactions and spectroscopy continues to accumulate from weak decays of kaons, charm (telling about the low-mass $I = J = 0$ dipion resonance $\sigma$), and $B$ (illuminating properties of scalar mesons like $f_0$ and $\alpha_0$, which must be understood if one is to identify glueballs), and radiative $\phi$ decays. For example, the NA48 Collaboration at CERN has obtained information on $\pi\pi$ scattering lengths from $K_{e4}$ and $K^+ \to \pi^+\pi^0\pi^0$ decays. Some results are summarized in Fig. 4.

Scattering lengths $a_J^I$ are conventionally labelled by total momentum $J$ and isospin $I$. The predictions of current algebra are $a_J^0 = -0.044$ and $a_J^0 = 0.22$. The NA48 measurement of $a_J^0$ seems to be slightly above this last value but more data from NA48 will tell whether there really is a discrepancy.
The helicity structure of $\rho$ mesons in the reaction $e^+e^- \rightarrow \rho^+\rho^-$ has recently been measured by the BaBar Collaboration, with the result $F_{00} = 0.54 \pm 0.10 \pm 0.02$, where the subscripts denote $\rho$ helicity. This is to be compared with the asymptotic prediction $F_{00} \rightarrow 1$. Should one be surprised? Are there related tests at comparable scales of $E_{CM} \approx 10$ GeV?

Recent results by the KLOE Collaboration shed light on the quark/gluon content of $\eta'$ through the decay $\phi \rightarrow \eta'\gamma$. Comparison of this decay with others (such as $\phi \rightarrow \eta\gamma$, $\rho \rightarrow \eta\gamma$, $\eta \rightarrow \gamma\gamma$, $\eta' \rightarrow \gamma\gamma$, and so on), following a method proposed some time ago, lead to the conclusion that the glue content of the $\eta'$ is $(14 \pm 4)\%$.

8 Charme

Results from BES were presented at this Conference on states reached in $J/\psi$ decays, including a broad $X(1580)$ decaying to $K^+K^-$ seen in $J/\psi \rightarrow K^+K^-\pi^0$ and an $\omega\phi$ threshold peak seen in $J/\psi \rightarrow \gamma\omega\phi$, as well as on multibody $\psi(2S)$ decays. CLEO results included confirmation of the $Y(2460)$ in a direct scan and in radiative return; a new measurement of $M(D^0)$ which implies that the $X(3872)$ is bound by $0.6\pm0.6$ MeV; and observation of $\psi''(3770) \rightarrow \gamma\chi_c$ decays with rates confirming its assignment as the $1^3D_1$ charmonium state. Belle reported two-photon production of several states including $Z(3930)$, a $\chi_c2(2P)$ candidate.

9 Charmed hadrons

9.1 $L = 0$ states

BaBar has identified the $\Omega_c^+$, a candidate for the lowest-lying $J = 3/2$ ccc state lying $70.8 \pm 1.0 \pm 1.1$ MeV above the $\Omega_c$ (also recently studied by BaBar). This mass splitting agrees with that predicted in the quark model. One now has a complete set of candidates for the $L = 0$ mesons and baryons containing a single charmed quark. As we shall see, charmed hadron masses are useful in anticipating those of hadrons containing a $b$ quark.

9.2 Orbitally-excited mesons

In the heavy-quark limit, mesons made of one heavy and one light quark are best described by coupling the light quark and the orbital angular momentum $L$ to a total $j$, and then $j$ to the heavy quark spin to form states of $J = j \pm 1/2$. For $L = 1$ one then has states with $j = 1/2$ (leading to $J = 0, 1$) and $j = 3/2$ (leading to $J = 1, 2$). The $J = 3/2$ states, predicted to be narrow, have been known for many years for both charmed-nonstrange and charmed-strange mesons. However, the $j = 1/2$ states, expected to be broad, proved more elusive.

The two $L = 1$, $j = 1/2$ $c\bar{s}$ mesons, the $D_{s0}(2317)$ and $D_{s1}(2460)$, were lighter than expected by most theorists. Lying below the respective $DK$ and $D^*K$ thresholds for strong decays, they turned out to be narrow, decaying radiatively or via isospin-violating $\pi^0$ emission. Their low masses were anticipated in schemes which pegged them as chiral partners of the $D_s$ and $D_s^*$. Regarding them as bound states of $DK$ and $D^*K$, respectively, they each would have a binding energy of 41 MeV. It would be interesting to see if a similar pattern holds for $B_{sJ}$ as $B^{(*)}K$ bound states. The lesson is that light-quark degrees of freedom appear to be important in understanding heavy-quark systems.

Higher-mass $c\bar{s}$ states have now been reported. The Belle Collaboration sees a $D_s$ state in the $M(DK^+)$ spectrum in $B^0 \rightarrow D_s^0D^0K^+$. It has $M = (2715 \pm 11^{+14}_{-15})$ MeV and $\Gamma = (115 \pm 20^{+36}_{-32})$ MeV. BaBar could be seeing this state, though not with significance. It has $J^P = 1^-$ and lies $603^{+16}_{-18}$ MeV above $D_s^*(2112)$, to be compared with $2S$--$1S$ splittings of $681 \pm 20$ MeV for $s\bar{s}$ and $589$ MeV for $c\bar{c}$. It appears to be a good $c\bar{s}(2S1)$ candidate.
Another $D_s$ state is seen decaying to $D^0 K^+$ and $D^+ K_S$.\textsuperscript{77} It has $M = (2856.6 \pm 1.5 \pm 5.0)\text{ MeV}$ and $\Gamma = (48 \pm 7 \pm 10)\text{ MeV}$. It can be interpreted as the first radial excitation of $D_{s0}(2317)$\textsuperscript{78} or a $J^P = 3^- (3D_2)$ state.\textsuperscript{79} Angular distributions of decay products should permit a distinction.

While the established (narrow) $j^P = 3/2^+$ states $D_1(2422)$, $D_2(2460)$ have been known for quite some time, there is more question about the broad $j^P = 1/2^+$ candidates. Both CLEO\textsuperscript{80} and Belle\textsuperscript{81} place the broad $j^P = 1/2^+$, $J^P = 1^+$ candidates in the narrow range 2420–2460 MeV, but Belle\textsuperscript{81} and FOCUS\textsuperscript{82} differ somewhat with respect to broad $j^P = 1/2^+$, $J^P = 0^+$ candidates, placing them only in a rather wide range 2300–2400 MeV.

One feature of note is that orbital excitation to the well-established $j = 3/2$ states costs (472,482) MeV for ($D_{s0}^*, D_0^*$). We shall compare this figure with a corresponding one for $B$ mesons.

10 Beauty hadrons

10.1 $L = 0$ states

CDF has observed $\Lambda_b$ and $\Sigma_b^0$ candidates decaying to $\pi^\pm \Lambda^0_b$.\textsuperscript{83,84} Their mass measurements are aided by a new precise value, also due to CDF,\textsuperscript{85} $M(\Lambda_b) = (5619.7 \pm 1.2 \pm 1.2)\text{ MeV}$. It is worth comparing this mass with a simple quark model prediction.

The light ($u,d$) quarks in $\Lambda_c$ and $\Lambda_b$ must be coupled to spin zero, by the requirements of Fermi statistics, as they are antisymmetric in color ($3^*$) and flavor ($I = 0$) and symmetric in space (S-wave). Aside from small binding effects, one then expects $M(\Lambda_b) - M(\Lambda_c) = M_b - M_c$, where $M_b$ and $M_c$ are “constituent” quark masses whose difference $M_b - M_c$ may be obtained from $B(\ell)$ and $D(\ell)$ mesons by taking the combinations $(3M^* + M)/4$ for which the hyperfine $Q\bar{q}$ interactions cancel. Using $[3M(B^*) + M(B)]/4 = 5314.6 \pm 0.5\text{ MeV}$ and $[3M(D^*) + M(D)]/4 = 1973.0 \pm 0.4\text{ MeV}$ one then finds $M_b - M_c = 3341.6 \pm 0.6\text{ MeV}$. (This is slightly larger than the difference between $M_b = 4796\text{ MeV}$ and $M_c = 1666\text{ MeV}$ reported by Kühn.\textsuperscript{86}) Combining this difference with $M(\Lambda_c) = 2286.46 \pm 0.14\text{ MeV}$, one then predicts $M(\Lambda_b) = 5628.1 \pm 0.7\text{ MeV}$, 8 MeV above the observed value. One could ascribe the small difference, which goes in the right direction, to reduced-mass effects. A similar exercise predicts $M(\Xi_c) \approx 5.8\text{ GeV}$ from $M(\Xi_c) = 2469\text{ MeV}$.

We now turn to the $\Sigma_{b}^{*(\pm)}$ states. The direct measurements are of $Q^{(\pm)} \equiv M(\Sigma_{b}^{*(\pm)}) - M(\pi^\pm) - M(\Lambda_b)$, and it is found (under the assumption $Q^{(*)} - Q^* = Q^+ - Q^-$, which is expected to be good to 0.4 MeV\textsuperscript{87}) that

$$Q^+ = 48.4 \pm 0.3 \pm 0.2\text{ MeV} , \quad Q^- = 55.9 \pm 1.0 \pm 0.2\text{ MeV} .$$

(7)

With the new CDF value of $M(\Lambda_b)$, these results then imply

$$M(\Sigma_{b}^-) = 5815.2 \pm 0.3 \pm 1.7\text{ MeV} , \quad M(\Sigma_{b}^+) = 5807.5 \pm 0.3 \pm 1.7\text{ MeV} ,$$

(8)

$$M(\Sigma_{b}^{*+}) = 5836.7 \pm 0.3 \pm 1.7\text{ MeV} , \quad M(\Sigma_{b}^{*-}) = 5829.0 \pm 0.3 \pm 1.7\text{ MeV} .$$

(9)

These masses are entirely consistent with quark model predictions. (See\textsuperscript{87} and references therein.) The $\Lambda$ hyperon may be denoted $[ud]s$, whereas $[ud]$ denotes a pair antisymmetric in flavor and spin, whereas the $\Sigma^{+,0,-}$ quark wavefunction may be written as $(\cdots)s$, with $(\cdots) = (uu),(ud),(dd)$ shorthand for a pair symmetric in flavor and spin. $S(\cdots) = 1$ then can couple with $S(s) = 1/2$ to give $J = 1/2 (\Sigma)$ or $3/2 (\Sigma^*)$, with hyperfine splitting $\propto 1/m_s$. The mass difference between the spin-1 and spin-0 diquarks, $M(\cdots) = [2M(\Sigma^*) + M(\Sigma)]/3 - M(\Lambda)$, can be calculated from the spin-weighted average of $M(\Sigma^*)$ and $M(\Sigma)$, in which hyperfine interactions cancel out. This result is the same calculated from baryons containing $s, c, or b$:

$$\frac{\Sigma + 2\Sigma^*}{3} - \Lambda = 205.1 \pm 0.3\text{ MeV} , \quad \frac{\Sigma_c + 2\Sigma_c^*}{3} - \Lambda_c = 210.0 \pm 0.5\text{ MeV} ,$$

(10)
Table 1: Candidates for $L = 1$, $J^P = 3/2^+$ $B$ mesons. Masses in MeV.

<table>
<thead>
<tr>
<th></th>
<th>$B_1$</th>
<th>$B_2$</th>
<th>$B_{s1}$</th>
<th>$B_{s2}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>CDF</td>
<td>5738.8±5.5±1</td>
<td>5739.2±3.5±2</td>
<td>5829.2±0.2±0.6</td>
<td>5839.6±0.4±0.5</td>
</tr>
<tr>
<td>D0</td>
<td>5720.8±2.5±5.3</td>
<td>5746±2.4±5.4</td>
<td>–</td>
<td>5839.1±1.4±1.5</td>
</tr>
</tbody>
</table>

The hyperfine splittings themselves also obey reasonable scaling laws. One expects $M(\Sigma^*) - M(\Sigma) \propto 1/m_s$ so splittings for charm and bottom should scale as $1/m_c$, $1/m_b$, respectively. The differences for $s, c,$ and $b$, 191.4 ± 0.4, 64.4 ± 0.8, 21.3 ± 2.0 MeV, are indeed approximately in the ratio of $1/m_c : 1/m_s : 1/m_b$.

10.2 $L = 1$ mesons

Results from CDF and D0, summarized by Filthaut, are shown in Table 1. Arguments similar to those for the $L = 0$ baryons in the previous subsection imply that one should expect $M(B_2) - M(B_1) \approx M(B_{s2}) - M(B_{s1}) \approx 13$ MeV. This pattern does not seem to emerge clearly from the data, which in any case give mixed signals regarding hyperfine splittings. One pattern which does seem fairly clear is that orbital $j = 3/2$ $B, B_s$ excitations cost $\sim 50$ MeV less than for $D, D_s$.

11 Importance of thresholds

Many hadrons discovered recently require that one understand nearby thresholds, a problem with a long history. As one example, the cross section for $e^+e^- \to \text{(hadrons)}$ has a sharp dip around a center-of-mass energy of 4.25 GeV, which is just below the threshold for the lowest-lying pair of charmed mesons ($D^0$ and $\bar{D}^0$) which can be produced in a relative $S$-wave. All lower-mass thresholds, such as $D\bar{D}$, $D\bar{D}^*$, and $D^*\bar{D}^*$, correspond to production in relative $P$-waves, so the corresponding channels do not open up as quickly. The $D^0\bar{D}^0$ (+ c.c.) channel is the expected decay of the puzzling charmonium state $Y(4260)$ if it is a hybrid ($c\bar{c}$ + gluon). But this channel is closed, so others (such as the observed $\pi\pi J/\psi$ channel) may be favored instead.

It is likely that the dip in $e^+e^- \to \text{(hadrons)}$ is correlated with a substantial suppression of charm production just before the $D^0\bar{D}^0$ channel opens up. The cross section for $e^+e^- \to D^*D^*$ (a major charm channel) indeed experiences a sharp dip at 4.25 GeV. Perhaps the peak $Y(4320) \to \pi^+\pi^-\psi(2S)$ seen by BaBar, with $M = 4324 \pm 24$ MeV, $\Gamma = 172 \pm 23$ MeV, is correlated with some other threshold.

Many other dips are correlated with thresholds [e.g., in the $\pi\pi$ $S$-wave near $2M(K)$ or $\gamma^* \to 6\pi$ near $2M(p)$]. The BaBar Collaboration recently has reported a structure in $e^+e^- \to \phi f_0(980)$ at 2175 MeV. It could be a hybrid $s\bar{s}$ candidate in the same way that $Y(4260)$ is a hybrid $c\bar{c}$ candidate. The assignment makes sense if $M_c - M_s \approx (M_Y - M_K)/2 = 1.04$ GeV.

12 Quark masses

J. H. Kühn has presented explicit formulae for the running of quark masses. High-order corrections to the Taylor series for the heavy quark vacuum polarization function $\Pi_Q(q^2)$ are a tour de force. One may expect interesting things from this group on high-order corrections.
to \( R = \sigma(e^+e^- \rightarrow \text{hadrons})/\sigma(e^+e^- \rightarrow \mu^+\mu^-) \). The moments \( M_n = \int dsR(s)/s^{n+1} \) give consistent masses, with \( m_c(m_c) = 1287 \pm 13 \text{ MeV} \) from \( n = 1 \) and \( m_b(m_b) = 4167 \pm 23 \text{ MeV} \) from \( n = 2 \). These results are an update of Ref.\(^{95}\). The pole masses \( M_b = 4796 \text{ MeV} \) and \( M_b = 1666 \text{ MeV} \) differ by 3130 MeV, a bit less than the phenomenological value of 3342 MeV mentioned earlier in the prediction of \( M(\Lambda_b) \). One caveat is that old CLEO data were used with an arbitrary renormalization. CLEO should come out soon with new \( R \) values below \( BB \) threshold but needs to present its data above \( BB \) threshold similarly. These data were taken in connection with a search for \( \Lambda_b \Lambda_b \) production.\(^{96}\)

A. Pineda has reminded me of a work\(^{97}\) in which \( \tilde{m}_b(\tilde{m}_b) = 4.19 \pm 0.06 \text{ GeV} \) is obtained from a non-relativistic sum rule. Kühn’s talk has a compilation of many other values. The uncertainty in \( m_c \), reduced by Kühn’s analysis, is an important part of the theoretical error in calculating \( B(b \rightarrow s\gamma) \).\(^{98}\)

Although the top quark mass has been measured with impressive accuracy (see below), it may be possible by studying threshold behavior in \( e^+e^- \rightarrow t\bar{t} \) to learn it to about 0.1 GeV.\(^{99}\)

13 Heavy flavor production

Calculations of hadronic charm production are in rough accord with experiment (though there remains some excess peaking for small azimuthal angle between charm and anticharm). While the description of beauty production has improved vastly in the past few years, there are still some kinematic regions where experiment exceeds theory.\(^{100}\) Incisive beauty–antibeauty correlation measurements still do not exist despite long-standing pleas.\(^{101}\) One looks forward to these at the LHC.\(^{102}\)

The quantitative understanding of quarkonium production still seems elusive. It demands soft gluon radiation, “adjustable” to the observed cross section. This is not the same as a first-principles calculation.

14 Fragmentation and jets

The correct description of fragmentation was a key ingredient in improving the agreement of \( b \) production predictions with experiment.\(^{100}\) At this conference new and/or upgraded Monte Carlo routines were reported.\(^{103,104}\) A useful detailed check of their hadronization features would be to compare their predicted multiplicities and particle particle species with CLEO data on hadronic \( \chi_c \) decays\(^{68}\) or hadronic bottomonium decays (which are being analyzed by CLEO). One could also imagine applying the global determination of fragmentation functions reported by Kumano\(^{105}\) to these questions.

Progress also has been reported with spinor-based multigluon methods;\(^{106,107}\) definition of \( b \)-jets;\(^{108}\) correction for the underlying event;\(^{109}\) exclusive \( p\bar{p} \rightarrow p\bar{p}X \) reactions;\(^{110}\) inclusive cross sections;\(^{111,112}\) and an infrared-safe-safe jet definition.\(^{113}\) Jets in heavy-ion collisions will be especially challenging.\(^{114}\)

15 \( W \) and top

New CDF values of \( (M_W = 80413 \pm 48) \text{ MeV} \) and \( \Gamma_W = (2032 \pm 71) \text{ MeV} \) have recently been reported.\(^{115}\) The new world averages, \( M_W = 80398 \pm 25 \text{ MeV} \) and \( \Gamma_W = (2095 \pm 47) \text{ MeV} \), are consistent with the Standard Model. In the latter there is very little room for deviations since no “oblique” \( (S,T) \) corrections are expected.\(^{116}\)

\[
\Gamma(W) = \frac{G_F M_W^3}{6\pi \sqrt{2}} \left[ 3 + 6 \left[ 1 + \frac{\alpha_s(M_W)}{\pi} \right] \right] = (2100 \pm 4) \text{ MeV}.
\]
Now information on top quark mass and production comes from CDF and D0.\textsuperscript{117} Examples of new measurements in the $\ell + \text{jets}$ channel are $m_t = (170.5 \pm 2.4 \pm 1.2) \text{ GeV}$ (D0) and $(170.9 \pm 2.2 \pm 1.40) \text{ GeV}$ (CDF). The present world average is now $m_t = (170.9 \pm 1.8) \text{ GeV}$, an error of 1.1%. This places further pressure on the Higgs mass. The Standard Model fit gives $M_H \leq 144 \text{ GeV}$ (95% c.l.), relaxed to 182 GeV if the present direct limit $M_H > 114.4 \text{ GeV}$ is considered.

One alternative to a light Higgs boson would involve custodial symmetry violation [for example, as provided by a new heavy SU(2) doublet with large mass splitting].\textsuperscript{118} Adding a vacuum expectation value ($V_0$) of a Higgs triplet with zero hypercharge which is only a few percent of the standard doublet $v = 246 \text{ GeV}$ would be sufficient to substantially relax the upper limit on $M_H$.\textsuperscript{119}

The D0 Collaboration sees single-top production at the expected level in three different analyses.\textsuperscript{2} CDF sees it in one analysis but not in two others.\textsuperscript{120} When the dust settles, this measurement is expected to provide useful information on $|V_{tb}|$.

16 Dibosons and Higgs

CDF and D0 have presented evidence for $WZ$ and $ZZ$ production, as summarized by F. Würthwein.\textsuperscript{121} D0 has seen a dip corresponding to the expected radiation zero in $W\gamma$ production. The subprocess $u\bar{d} \rightarrow W^+\gamma$ has a zero at $\cos\theta_{CM} = -1/3$, while $\bar{u}d \rightarrow W^-\gamma$ has a zero at $\cos\theta_{CM} = 1/3$.

In a search for the Higgs boson in the $H \rightarrow \tau\tau$ channel, bounds from CDF are “degraded” thanks to an excess of events for $M_H \simeq 160 \text{ GeV}$. On the other hand, D0 sees a deficit there.\textsuperscript{122} This mass range may be the first interval accessible with 8 fb$^{-1}$ at the Tevatron; sensitivities are improving faster than $1/\sqrt{L dt}$.\textsuperscript{123} It would be wonderful if a way were found to extend the run!

An interesting scheme for generating the Higgs boson via spontaneous conformal symmetry breaking was presented.\textsuperscript{124} As this tends to give a fairly heavy Higgs boson, it must be confronted with the tightening precision electroweak constraints. Strong electroweak symmetry breaking scenarios also were described.\textsuperscript{125} These essentially adapt chiral models to the TeV scale, replaying the strong interactions at a factor $v/f_\pi \approx 2650$ higher in energy. Light-Higgs scenarios are not ruled out; for instance, it has been asked whether the mass of the $b\bar{b}(1^1S_0)$ state, the as-yet-unseen $\tau_b$, is standard or is affected by mixing with a light Higgs boson.\textsuperscript{126} One Standard Model prediction\textsuperscript{127,128} is $M(\tau_b) = 9421 \text{ MeV}$.

Higgs decays to multiparticle final states have been described using twistor methods.\textsuperscript{107} It may be possible to produce a Higgs boson at LHC in the double-diffractive reaction $pp \rightarrow ppH$, monitoring the small-angle protons using Roman pots.\textsuperscript{110} One problem will be distinguishing which of the multiple interactions per crossing was the source of the scattered protons. This pileup effect may be solvable if one can make sufficiently rapid trigger decisions.

Two-Higgs models, if confirmed, provide a gateway to supersymmetry.\textsuperscript{129} Such proliferation of the Higgs spectrum, entailing two charged and three neutral Higgs bosons, also is a feature of grand unified theories beyond the minimal SU(5), such as SO(10).

17 Proton structure and diffraction

The proton spin $1/2$ is composed of $\frac{1}{2}\Delta \Sigma + \Delta G + \Delta L$, corresponding respectively to quarks, gluons, and orbital angular momentum. $\Delta \Sigma \simeq 0.3$; what’s the rest? The COMPASS$^{130}$ and STAR$^{131}$ Collaborations have shown that $\Delta G$ is not enough; one must have $\Delta L > 0$.\textsuperscript{\textit{End}}
Neutral-current $e\nu$ interactions at HERA have displayed the first evidence for parity violation in high-$Q^2$ deep inelastic scattering. HERA is helping to pin down structure functions and their evolution for use at the LHC. Also at HERA, it has been found that the Pomeron slope is different in $\rho^0$ and $J/\psi$ photoproduction. These reactions correspond respectively to soft and hard processes.

18 Heavy ion collisions

One has seen the adaptation of string theory ideas to properties of the quark-gluon plasma: hydrodynamic properties involve previously intractable strong-coupling calculations. In heavy-ion jet production, the recoiling jet is quenched if it must pass through the whole nucleus. This provides information about the properties of nuclear matter. An interesting rapidity “ridge” is seen in many processes. Could this be a manifestation of QCD “synchrotron radiation”? Do previous emulsion experiments display this feature?

One way to describe nuclear matter effects is via medium-modified fragmentation functions probe nuclear matter effects. Useful information is provided by $\gamma-\pi^0$ and $\gamma-\gamma$ correlations. Hanbury-Brown-Twiss correlations between identical particles (e.g., $\pi^+\pi^-$) provide information on the viscosity of the quark-gluon plasma and on the geometry and time evolution of the “hot” region. Charmed particles are found to interact with the nuclear medium in the same way as others. It is not clear whether there is a difference between the interactions of $c\bar{c}$ and $\bar{c}c$ states; certainly $K^+$ and $K^-$ do interact differently with nonstrange matter. Other important issues in nuclei include low-$x$ parton saturation and the question of whether quarkonium suppression is taking place.

19 Beyond the Standard Model

As this is a large field, I would like to comment on just a few items which I consider especially worth watching in the next few years.

1. The muon's $g-2$ value can get big contributions in some SUSY models. In units of $10^{-11}$, $a_\mu \equiv (g_\mu-2)/2 = 116 591 793$ (68) (theory), to be compared with 116 592 080 (63) (experiment). These differ by $(287 \pm 93)$ or $3.1\sigma$. This relies upon evaluating hadronic vacuum polarization via $e^+e^-$ annihilation. If one uses $\tau$ decays the discrepancy drops to $1.2\sigma$. The inconsistency is worth sorting out.

2. Non-standard explanations abound for the deviation of the effective $\sin(2\beta)$ in $b \rightarrow s$ penguins from the “reference value” obtained in decays dominated by $b \rightarrow c\bar{c}s$. The current biggest discrepancy is in $S_{K_S} = 0.33 \pm 0.21$, versus a nominal value of $0.678 \pm 0.026$. This could be due, for instance, to exchange of a new $Z'$ masquerading as an electroweak penguin. The study of $b \rightarrow s\ell^+\ell^-$ and searches at the Tevatron and LHC will see or bound $Z'$ effects. Forward-backward asymmetries can be quite sensitive to $Z'^*$s. One will be able to study such asymmetries at the LHC by passing to non-zero pseudorapidity $\eta$.

The $b \rightarrow s\ell^+\ell^-$ decays show no anomalous behavior so far. Belle/BaBar differ a bit and CDF agrees with BaBar with $B(B^0 \rightarrow K^0\mu^+\mu^-) = (0.82 \pm 0.31 \pm 0.10) \times 10^{-6}$, and with Belle with $B(B^+ \rightarrow K^+\mu^+\mu^-) = (0.60 \pm 0.15 \pm 0.04) \times 10^{-6}$.

3. It is encouraging to see the results searches for a right-handed $W$: $M_{W_R} > (790, 760)$ MeV for $M_{W_R}(\nu_R, \nu_{\tau})$ $M_{\nu_R}$. The case of a right-handed $\nu_R$ heavier than $M_{W_R}$, in particular, means that one must search for $W_R$ in the hadronic channel $tb$. 

120, 148 $M_{W_R}$
20  Dark matter in many forms

Ordinary matter exists in several stable forms: $p$, $n$ (when incorporated into nuclei), $e^-$, three flavors of neutrinos $\nu$ [$\nu_{\mu, \tau} \gg \nu_{\text{Universe}}$]. We could expect dark matter (5–6 x ordinary matter) to exhibit at least as much variety, for example if its quantum numbers are associated with a big gauge group largely shielded from current observations. \textsuperscript{150} “Mirror particles,” reviewed extensively by Okun, \textsuperscript{151} are one example of this possibility.

There are at least two well-motivated dark matter candidates already (axions and neutralinos). Axion dark matter has not received the attention it deserves. RF cavity searches are going slowly; there is a large range of frequencies still to be scanned with enough sensitivity. Some variants of supersymmetry have long-lived next-to-lightest superpartners, decaying to the lightest superpartners over a detectable distance. Charged and neutral quasi-stable candidates\textsuperscript{152} could be split by so little that they charge-exchange with the detector, implying new tracking signatures.

Dark matter could have non-zero charges purely in a hidden sector and thus be invisible to all but gravitational probes. Such opportunities might be provided by the LISA detector. \textsuperscript{153}

Experience with hadron physics may help us deal with unexpected dark matter forms and interactions. This could be so, for example, if investigations at the TeV scale uncover a new strongly-interacting sector, as expected in some theories of dynamical electroweak symmetry breaking.

21  Outlook

Impressive measurements from BaBar, Belle, CDF, CERN NA48, CLEO, D0, KLOE, RHIC, and other experiments have provided much fuel for theoretical interpretations at this conference. The understanding of hadron physics plays a key role. Much knowledge about fundamental electroweak interactions relies on separating out the strong interactions. Methods include theoretical calculations (pQCD, SCET) and correlation of measurements through flavor symmetry. Conversely, low-energy hadron physics has benefitted greatly from weak interactions; $K$, $D$, $B$ decays have provided information on $\pi\pi$ scattering, $\sigma$ and other scalar mesons, and patterns of final-state interactions which go beyond what perturbative methods can anticipate.

Experiments at the Tevatron have shown that one can do excellent flavor physics in a hadronic environment. We look forward to fruitful results from LHCb on $B_s \to \mu\mu$, CP violation in $B_d \to J/\psi\phi$, and many other topics.

Higgs boson searches are gaining in both sensitivity and breadth; gaps are being plugged. In addition to the discovery of the Higgs at the LHC (unless Fermilab finds it first!), we can look forward to measurements of $\sigma_T$, flavor, top, Higgs, new particles and forces.

Discussions of a super-B-factory, possibly near Frascati, are maturing.\textsuperscript{154} Such a machine might solve the $b \to s$ penguin problem once and for all. With a luminosity approaching 100 times current values, it would permit tagging with fully reconstructed $B$'s in all those final states not studied with partial tags. Upgrades of KEK-B and LHCb also are being contemplated. Finally, neutrino studies\textsuperscript{155} (near-term and more ambitious) and the ILC are also on our horizon. Our field has much to look forward to in the coming decades.

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102. V. Andreev, this Conference.
103. F. Krauss, this Conference.
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106. D. Varman, this Conference.
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112. O. Norniella, this Conference.
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114. J. Putschke and C. Salgado, talks at this Conference.
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121. F. Würtzwein, this Conference.
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125. S. Allwood, this Conference.
128. A. Pineda, this Conference.
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131. K. Kowalik, this Conference.
132. S. U. Noor, this Conference.
133. T. Danielson, this Conference.
134. T. Hreus, this Conference.
135. J. Casalderrey-Solana, this Conference.
137. L. Cunqueiro, this Conference.
138. Z. Belghebsi, this Conference.
139. P. Romatschke and P. Chaloupka, talks at this Conference. See also P. Romatschke, arXiv:nucl-th/0701032.
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152. A. Rizzi, this Conference.
155. B. Kayser, this Conference.
# Rencontres de Moriond 2007

## QCD and High Energy Hadronic Interactions

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