MOMENTUM CORRELATIONS IN $Z^0 \rightarrow b\bar{b}$ AND THE MEASUREMENT OF $R^0_b$

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Abstract

We study the correlations of $b\bar{b}$ pairs produced in inclusive $Z^0$ decays, and their possible influence upon the determination of $R_b$.

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1 Introduction

The measured value of $R^0_b = \Gamma_{b\bar{b}}/\Gamma_{\text{had}}$ in $Z^0$ decays is drawing considerable attention in the physics community. On the one hand, according to all four LEP experiments, it deviates significantly from the Standard Model prediction, a result that has remained quite stable in the last few years [1]–[7]. On the other hand, it is a considerably difficult measurement, and it is certainly disturbing that the only deviation from the Standard Model is found in a quantity that is affected by complicated strong-interaction effects. In the present paper we examine one possible source of systematic errors that may affect the measurement of $R^0_b$, whose origin is purely dynamical. In several experimental techniques for the measurement of $R^0_b$, the tagging efficiency is extracted from data by comparing the sample of events in which only one $b$ has been tagged, with the one in which both $b$’s are observed. If the production characteristics of the $b$ and the $\bar{b}$ were completely uncorrelated, this method would yield the exact answer with no need for corrections. Of course, other correlations of an experimental nature should be properly accounted for, but their discussion is outside the scope of the present theoretical paper. Standard QCD gluon emission generates a correlation of the quark–antiquark momenta of the order of $\alpha_s$. Other dynamical effects, such as the production of heavy quark pairs via a gluon-splitting mechanism, may affect the measurement. However, they are to some extent better understood. In the first place they tend to give soft heavy quarks, and they are therefore easily eliminated. Furthermore, experimental studies of these production mechanisms have begun to appear [8].

The aim of the present work is to study the dynamical correlation of the heavy flavoured quark–antiquark pair, with particular emphasis upon its impact on the determination of $R_b$. We will use leading-order QCD formulae throughout, emphasizing the cases in which a higher-order calculation would be desirable and possible. We found that a particularly relevant quantity is the average momentum correlation $r$ defined as

$$r = \frac{\langle x_1 x_2 \rangle - \langle x \rangle^2}{\langle x \rangle^2}, \tag{1.1}$$

where $x_{1(2)}$ are the Feynman $x$ of the produced $B(\bar{B})$ mesons, $x_{1(2)} = p_{B(\bar{B})}/p^{(\text{max})}_B$, and $\langle x \rangle = \langle x_{1(2)} \rangle$. In fact, we will show that under the assumption that the detection efficiency for a $B$ meson is nearly proportional to its momentum, the quantity $r$ represents the relative correction to $R_b$ due to momentum correlations. The quantity
$r$ has a well-behaved perturbative expansion, starting at order $\alpha_s$, in spite of the fact that the perturbative expansions of $\langle x_1 x_2 \rangle$ and $\langle x \rangle$ are not well-behaved, since their coefficients are enhanced by logarithms of the total annihilation energy over the heavy-quark mass. Our perturbative calculation for $r$ gives a value of the order of 1%. Power-suppressed corrections may still affect this quantity to some extent. If power corrections of the order of $1/Q$ are present, they may turn out to be of the same order of the perturbative value at LEP energy. Monte Carlo models \cite{9} seem to support this possibility \cite{10}.

In general, the QCD formalism allows us to compute the double differential cross section for the inclusive production of the heavy-flavour pair, using as input the value of $\alpha_s$ and the fragmentation function of heavy-flavoured mesons. It is easy to compute such quantities in leading order. The leading order result may be improved on one hand by including next to leading corrections, and by resumming terms that are enhanced logarithmically near the end of the phase space (i.e. for $x_{1(2)} \to 1$). At present these calculations are not available.

Shower Monte Carlo models include to some extent perturbative effects of leading order, and effects that are enhanced near the end of the phase space. At the parton level, they should therefore provide a description that is compatible with the leading-order calculations reported in the present paper. Some recent studies \cite{10} support this conclusion. On the other hand, fixed order matrix elements Monte Carlo programs should be used cautiously when computing correlations, because of problems arising in the truncation of the perturbative expansion.

The rest of the paper is organized as follows. In Section 2 we discuss the average momentum correlation, and how it affects the determination of $R_b$. In Section 3 we describe the leading-order calculation of the double inclusive cross section for heavy flavour production. Finally, we give our conclusions in Section 4.

2 The average momentum correlation

We begin by considering the simple case in which the efficiency for tagging a $B$ meson is a linear function of its momentum. This simplifying assumption allows us to make very precise statements about the correlation. Furthermore, it is not extremely far from reality, in the sense that the experimental tagging efficiency is often a growing
function of the $B$ momentum. Let us therefore assume

$$\epsilon_b = Cx_1, \quad \epsilon_{\bar{b}} = Cx_2,$$

where $\epsilon_b$ is the tagging efficiency. We assume that in our ideal detector the detection efficiency is not influenced by the presence of another tag. Therefore the total number of tags is then given by

$$\frac{N_1}{N} = R_0^0 C (\langle x_1 \rangle + \langle x_2 \rangle) = R_0^0 2C \langle x \rangle,$$

where $x_2$ is the Feynman $x$ of the produced $\bar{B}$ meson, and with $\langle x \rangle$ we indicate the common value of their average. Observe that the in the above equation $N_1$ is the number of tags, given by the number of events with one tag plus twice the number of events with two tags. On the other hand, the number of events with two tags will be given by

$$\frac{N_2}{N} = R_0^0 C^2 \langle x_1 x_2 \rangle = R_0^0 (C \langle x \rangle)^2 \times (1 + r),$$

where $r$ is defined in eq. (1.1). We can then extract $R_0^0$

$$R_0^0 = \frac{N_1^2}{4NN_2} \times (1 + r).$$

The quantity $r$ cannot be measured; one therefore has to compute it in order to determine $R_0^0$.

The basic QCD processes for the production of a $b\bar{b}$ pair are depicted in fig. 1. Let us focus on the process up to the order $\alpha_s$. The final state is fully defined in this case by the variables $\bar{x}_1 = 2\epsilon_b/E$ and $\bar{x}_2 = 2\epsilon_{\bar{b}}/E$, where $E$ is the total centre-of-mass energy. The gluon momentum fraction $x_3$ is given by $2 - \bar{x}_1 - \bar{x}_2$. We can write the differential cross section, normalized to 1, as

$$\frac{d\sigma}{d\bar{x}_1d\bar{x}_2} = \delta(1 - \bar{x}_1)\delta(1 - \bar{x}_2) + \alpha_s \frac{d\sigma^{(1)}}{d\bar{x}_1d\bar{x}_2}.$$
The term of order $\alpha_s$ is singular when $\bar{x}_1$ and $\bar{x}_2$ approach 1. In fact, in this limit the emitted gluon is soft, and one expects a $1/x_3$ singularity. Furthermore, virtual corrections, concentrated in the region $\bar{x}_1 = 1$ and $\bar{x}_2 = 1$ are also infrared-divergent. Since our differential cross section is normalized to 1, it is easy to write down a formula for the average of some physical quantity $G(\bar{x}_1, \bar{x}_2)$ without having to explicitly introduce the virtual corrections. We simply write

$$\langle G(\bar{x}_1, \bar{x}_2) \rangle = G(1, 1) + \alpha_s \int d\bar{x}_1 d\bar{x}_2 \frac{d\sigma^{(1)}}{d\bar{x}_1 d\bar{x}_2} (G(\bar{x}_1, \bar{x}_2) - G(1, 1)).$$

(2.6)

The integral is extended in the appropriate Dalitz region. The subtraction term under the integral sign embodies the effects of the virtual corrections, and of the normalization to the total cross section. Thus, if we choose $G(\bar{x}_1, \bar{x}_2) = 1$ we clearly get 1, which is our normalization. Appropriate formulae for $d\sigma^{(1)}/d\bar{x}_1 d\bar{x}_2$ can be found for example in ref. [11]. Formula (2.6) can be easily implemented in a computer program. We computed the following quantities

$$c_1 = \int d\bar{x}_1 d\bar{x}_2 \frac{d\sigma^{(1)}}{d\bar{x}_1 d\bar{x}_2} (x_1 - 1),$$

(2.7)

$$d_1 = \int d\bar{x}_1 d\bar{x}_2 \frac{d\sigma^{(1)}}{d\bar{x}_1 d\bar{x}_2} (x_1 x_2 - 1),$$

(2.8)

$$e_1 = \int d\bar{x}_1 d\bar{x}_2 \frac{d\sigma^{(1)}}{d\bar{x}_1 d\bar{x}_2} (x_1 x_2 \text{cut}(x_1, x_2) - 1),$$

(2.9)

where

$$\text{cut}(x_1, x_2) = \theta(\max(x_1, x_2) - x_3).$$

(2.10)

The theta function in the last equation corresponds to the requirement that in doubly tagged events the $b$ and the $\bar{b}$ are in opposite hemispheres with respect to the thrust axis. We have

$$\langle x \rangle = 1 + \alpha_s c_1, \quad \langle x_1 x_2 \rangle = 1 + \alpha_s d_1, \quad \langle x_1 x_2 \text{cut}(x_1, x_2) \rangle = 1 + \alpha_s e_1.$$  

(2.11)

We find, for $m_b = 5$ GeV,

$$c_1 = -1.142, \quad d_1 = -2.185, \quad e_1 = -2.223.$$  

(2.12)

Observe that the coefficients we found are quite large, since one would have naively expected them to be of the order of $1/\pi$. This is because large logarithms $\log E/m_b$
arise in the coefficients, thus making the perturbative expansion for the above quantities quite unreliable. The quantities \( r \) and \( r' \), however, have small perturbative coefficients

\[
r = \frac{\langle x_1 x_2 \rangle - \langle x \rangle^2}{\langle x \rangle^2} = \alpha_s (d_1 - 2c_1) = 0.099 \alpha_s
\]

\[
r' = \frac{\langle x_1 x_2 \text{cut}(x_1, x_2) \rangle - \langle x \rangle^2}{\langle x \rangle^2} = \alpha_s (e_1 - 2c_1) = 0.061 \alpha_s.
\] (2.13)

Observe that the coefficients are even below the expected magnitude of \( 1/\pi \). A better illustration of the presence of large logarithms, and of their cancellation in the quantities \( r \) and \( r' \) is given in table 1, where the coefficients \( c_1 \), \( d_1 \) and \( e_1 \) have been computed for \( m = 10, 1, \) and \( 0.1 \) GeV. While \( c_1 \), \( d_1 \) and \( e_1 \) grow as the mass decreases, \( r \) and \( r' \) approach a finite limit. This is due to the fact that the large logarithms present in \( e_1 \), \( d_1 \) and \( e_1 \) cancel in the linear combinations appearing in \( r \). It is easy to prove that this cancellation must occur to all orders in perturbation theory. In fact, according to the factorization theorem, the double inclusive cross section \( b \bar{b} \) production, in the limit of \( E \gg m_b \), can be written as

\[
\frac{d\sigma}{dx_1 dx_2} = \int \frac{d\hat{\sigma}}{dy_1 dy_2} D(z_1) D(z_2) \delta(y_1 z_1 - x_1) \delta(y_2 z_2 - x_2) dy_1 dz_1 dy_2 dz_2, \quad (2.14)
\]

where (in the limit \( E/m_b \to \infty \)) \( \hat{\sigma} \) has a perturbative expansion in \( \alpha_s \) with finite coefficients (observe that the distinction between the momentum- or energy-defined Feynman \( x \) becomes irrelevant in the limit we are considering here). The divergent terms are all absorbed in the fragmentation functions \( D(z) \). We then have

\[
\langle x \rangle = \int dx_1 dx_2 \ x_1 \frac{d\sigma}{dx_1 dx_2} = \left( \int dz \ zD(z) \right) \int dy_1 dy_2 \ y_1 \frac{d\hat{\sigma}}{dy_1 dy_2}, \quad (2.15)
\]

\[
\langle x_1 x_2 \rangle = \int dx_1 dx_2 \ x_1 x_2 \frac{d\sigma}{dx_1 dx_2} = \left( \int dz \ zD(z) \right)^2 \int dy_1 dy_2 \ y_1 y_2 \frac{d\hat{\sigma}}{dy_1 dy_2}. \quad (2.16)
\]

| \( m \) | \( c_1 \) | \( d_1 \) | \( e_1 \) | \( r \) | \( r' \) |
|-----|-----|-----|-----|-----|-----|
| 10 GeV | -0.773 | -1.460 | -1.486 | 0.086 \( \alpha_s \) | 0.060 \( \alpha_s \) |
| 1 GeV | -2.053 | -4.000 | -4.045 | 0.106 \( \alpha_s \) | 0.061 \( \alpha_s \) |
| 0.1 GeV | -3.348 | -6.590 | -6.636 | 0.106 \( \alpha_s \) | 0.060 \( \alpha_s \) |

Table 1: Mass dependence of the coefficients \( c_1 \), \( d_1 \) and \( e_1 \).
In the ratio $\langle x_1 x_2 \rangle / \langle x \rangle^2$ (and therefore in $r$) the integral containing $D$ cancels. Thus, the perturbative coefficients of $r$ are finite in the limit $E/m_b \to \infty$. Observe that in the derivation of eq. (2.14) we have assumed the relation
\[
\int D(z) \, dz = 1,
\]
which is appropriate when we can neglect the secondary production of $b\bar{b}$ pairs via gluon splitting.

Since the cancellation of large logarithms takes place order by order in perturbation theory, it would be wrong to include incomplete higher-order corrections in the expression for $r$. Thus, for example, if we write
\[
r = \frac{1 + d_1 \alpha_s - (1 + c_1 \alpha_s)^2}{1 + c_1 \alpha_s} = (d_1 - 2c_1)\alpha_s + c_1(3c_1 - 2d_1)\alpha_s^2 + \cdots ,
\]
the coefficient in front of the term of order $\alpha_s^2$ is very large when compared to the coefficient in front of the term of order $\alpha_s$. This is due to the fact that there are missing terms in the $\mathcal{O}(\alpha_s^2)$. If a complete $\mathcal{O}(\alpha_s^2)$ calculation is performed, these large terms cancel. Although a calculation of the next-to-leading term is in principle possible, it is not yet available. A calculation of the process of $b\bar{b}gg$ production has been given in ref. [13], but the virtual corrections for the $b\bar{b}g$ process in the case of massive quarks are not available. In order to give an estimate of the next-to-leading effects, and to verify the cancellation of the large terms in the $\mathcal{O}(\alpha_s^3)$ coefficients, we have computed the $\mathcal{O}(\alpha_s^2)$ contribution to $r$ coming from the real emission process $b\bar{b}gg$ alone. We would like to emphasise that these results can be considered at most as an estimate of the order of magnitude of the $\mathcal{O}(\alpha_s^3)$ coefficient, since the virtual terms are missing. Appropriate subtraction terms have been included in order to regulate the soft singularities. We get
\[
c_2 = -2.845, \quad d_2 = -4.595, \quad e_2 = -4.499.
\]
We have
\[
r = (d_1 - 2c_1)\alpha_s + (d_2 - 2c_1d_1 - 2c_2 + 3c_1^2)\alpha_s^2 + \mathcal{O}(\alpha_s^3)
\]
\[
= 0.099\alpha_s + 0.017\alpha_s^2 + \mathcal{O}(\alpha_s^3),
\]
\[
r' = (e_1 - 2c_1)\alpha_s + (e_2 - 2c_1e_1 - 2c_2 + 3c_1^2)\alpha_s^2 + \mathcal{O}(\alpha_s^3)
\]
\[
= 0.061\alpha_s + 0.026\alpha_s^2 + \mathcal{O}(\alpha_s^3),
\]
(2.18)
which shows a well-behaved perturbative expansion. Notice that if we do not include the $c_2$, $d_2$ and $e_2$ terms, the coefficients of the term of order $\alpha_s^2$ are an order of magnitude larger than the coefficient of the term of order $\alpha_s$, which again shows the importance of including all enhanced terms.

The correlation we computed is a small effect, of the order of 1%. In fact, we have to decide what is the scale at which $\alpha_s$ ought to be evaluated in eqs. (2.13). In view of our discussion on the cancellation of large logarithms, the correlation should be dominated by large momenta, and therefore the appropriate scale should be of the order of $E$. In fact, experience with jet physics suggests the use of a somewhat smaller value. For example, the range $0.12 < \alpha_s < 0.16$ gives a value of $r'$ between 0.7 and 1%.

3 Double inclusive cross section

Up to now, we have assumed in our discussion that the efficiency is linear in the $B$ momentum. Even in the more realistic case, in which the efficiency is a more complicated function of the kinematic variables, it is possible to compute the inclusive cross section for the production of a $b\bar{b}$ pair, provided one also knows the $B$ fragmentation function, which is to some extent measured at LEP. The appropriate formula is given in eq. (2.14). The expression for the short-distance cross section $\hat{\sigma}$ up to the order $\alpha_s$ is

$$
\frac{d\hat{\sigma}}{dy_1 dy_2} = \delta(1-y_1)\delta(1-y_2) + \frac{2\alpha_s}{3\pi} \frac{y_1^2 + y_2^2}{(1-y_1)(1-y_2)} .
$$

The + distribution sign specifies the way that the singularities at $y_1 = 1$ and $y_2 = 1$ should be treated. For any smooth function of $y_1$, $y_2$ we define

$$
\int_{y_1+y_2>1} dy_1 dy_2 \left[ \frac{y_1^2 + y_2^2}{(1-y_1)(1-y_2)} \right]_+ G(y_1, y_2) =
$$

$$
\int_{y_1+y_2>1} dy_1 dy_2 \left[ \frac{y_1^2 + y_2^2}{(1-y_1)(1-y_2)} \right]_+ (G(y_1, y_2) - G(1, y_2) - G(y_1, 1) + G(1, 1)) .
$$

Eqs. (2.14) and (3.1) consistently gives

$$
\frac{d\sigma}{dx_1} = \int dx_2 \frac{d\sigma}{dx_1 dx_2} = D(x_1) .
$$

The above equation fixes the factorization scheme to be the annihilation scheme as defined in ref. [12]. Observe that, in this case, the choice of scheme, i.e. eq. (3.3),

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As an assistant, I've transcribed the text from the image. It discusses the importance of including all enhanced terms in perturbative expansions of a well-behaved expansion. The correlation computed is a small effect, but the scale at which $\alpha_s$ should be evaluated is discussed. The double inclusive cross section is computed in a more realistic case, considering efficiencies as functions of kinematic variables. Eqs. 2.14 and 3.1 provide the formula for the short-distance cross section up to order $\alpha_s$, and the choice of factorization scheme is fixed to the annihilation scheme.
fixed unambiguously the result without the need of computing explicitly the virtual corrections. In fact, the most general formula for the short distance cross section is obtained by adding to eq. (3.1) terms of the form

$$\frac{\alpha_s}{\pi} \left[ \delta(1 - y_1)f(y_2) + f(y_1)\delta(1 - y_2) + g\delta(1 - y_1)\delta(1 - y_2) \right], \quad (3.4)$$

where $f$ is a generic function (in general a distribution) of one variable and $g$ is a number. However, in order for eq. (3.3) to be respected, these terms must be absent. From eqs. (2.14), (3.1) and (3.2) we immediately derive the $O(\alpha_s)$ formula

$$\frac{d\sigma}{dx_1 dx_2} = D(x_1)D(x_2) + \frac{2\alpha_s}{3\pi} \int_0^1 dy_1 dy_2 \theta(y_1 + y_2 - 1) \frac{y_1^2 + y_2^2}{(1 - y_1)(1 - y_2)} \times$$

$$D\left(\frac{x_1}{y_1}\right)D\left(\frac{x_2}{y_2}\right) \frac{1}{y_1y_2} \theta(y_1 - x_1) \theta(y_2 - x_2) - D(x_1)D\left(\frac{x_2}{y_2}\right) \frac{1}{y_2} \theta(y_2 - x_2)$$

$$- D\left(\frac{x_1}{y_1}\right)D(x_2) \frac{1}{y_1} \theta(y_1 - x_1) + D(x_1)D(x_2) \right]. \quad (3.5)$$

As an illustration, we plot in fig. 2 the double inclusive cross section $d\sigma/dx_1 dx_2$ as a function of $x_1$ for several values of $x_2$. We use the Peterson parametrization of the fragmentation function

$$D(z) = N_\epsilon \frac{z(1 - z)^2}{(1 - z)^2 + \epsilon^2}, \quad (3.6)$$

where $N_\epsilon$ is fixed by the condition $\int D(x)dx = 1$. We used the value $\epsilon = 0.04$, which gives $\langle x \rangle = 0.70$, and $\alpha_s = 0.12$. The positive momentum correlation is quite visible in the figure. As $x_2$ increases, the peak of the distribution in $x_1$ also moves towards larger values. Using the above formula, we can again compute $r$. The result should not depend upon the choice of the fragmentation function. In this case, the quantity $\langle x \rangle$ does not receive corrections of order $\alpha_s$. We get

$$\langle x \rangle = 0.7036, \quad \langle x_1x_2 \rangle = 0.4950 + 0.0525 \alpha_s, \quad r = 0.1061 \alpha_s, \quad (3.7)$$

which is close to the value previously obtained $r = 0.099$. In fact, the difference is due to mass effects, as can be seen from table [4].

There are several ways in which the above leading order calculation could be improved in principle. First of all, in the region of $x_1$, $x_2$ near 1, large logarithms of $1 - x$ arise, and they could be resummed to all order in perturbation theory. It would also be desirable to include next-to-leading corrections to the partonic cross section. At this moment, these corrections have not yet been computed.
4 Conclusion

Monte Carlo programs implement to some extent some of the QCD dynamics discussed in the present work. We therefore expect that, in general, at the parton level, they should give results consistent with the QCD calculation reported here, at least for quantities like $r$, in which large logarithmic effects cancel. Observe that the computation of $r$ must be performed either consistently at some fixed order, or including all logarithmically enhanced terms. Failure to do so may expose large, uncancelled logarithmic terms, and it may therefore lead to a wrong result. Shower Monte Carlo programs include all logarithmically enhanced terms. On the other hand, matrix element Monte Carlo programs use a truncated perturbative expansion. They can be used to compute $r$ only if each fixed-order contribution can be isolated, and all terms of order higher than the accuracy of the Monte Carlo can be thrown away.

The QCD factorization theorem guarantees that in quantities like $r$ large logarithmic effects, as well as hadronization effects, should cancel. This statement is valid in a leading-twist sense. Thus, it is possible that power-suppressed corrections to $r$ are present. A correction of the order of $1/Q$ would be a 1% effect, comparable to
the $\mathcal{O}(\alpha_s)$ corrections. It is difficult to say anything about the importance of power corrections, since our knowledge of the hadronization mechanism is quite poor. One has therefore to rely upon hadronization models, which are commonly implemented in Monte Carlo programs. The formalism presented in this paper can also be applied to the computation of other processes, such as, for example, the correlation in the inclusive production of strange mesons in the large-$x$ region. These processes may provide a testing ground for the theoretical predictions, and may help us gain confidence on the reliability of the hadronization models implemented in the Monte Carlo programs.

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