Measurement of the b-jet cross section at Vs=1.96 TeV
Peters, O.

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Chapter 1

The Standard Model

The Standard Model is the theory that describes (sub-)nuclear matter and its interactions. It encompasses the electromagnetic force and weak forces, combined in the electroweak theory, and the strong force described by the theory of quantum chromodynamics, QCD. The model is a Quantum Field theory based on the concept of local gauge invariance. For the Standard Model, the gauge symmetry group is $SU(3) \times SU(2)_L \times U(1)$, where $SU(3)$ is the symmetry group describing the strong interaction and $SU(2)_L \times U(1)$ the symmetry group describing the electroweak interaction. The theory describes two general classes of particles: spin-$\frac{1}{2}$ matter particles known as fermions and spin-1 gauge vector particles known as bosons, which carry the forces. In this formalism, $SU(2)_L$ involves only left-handed fermions (hence the $L$ subscript). The fermions can be subdivided in leptons and quarks, which can be grouped into three families. Some properties of the leptons and quarks are summarized in tables 1.1 and 1.2, respectively [22]. The vector gauge bosons that mediate the forces between these particles are required by the local gauge invariance to be massless. The bosons that carry the electroweak force - namely, two charged particles, $W^+$ and $W^-$ and one neutral particle $Z^0$ - interact with themselves through the triple gauge couplings. Through the Higgs mechanism, they acquire a non-zero mass by spontaneous symmetry breaking of the $SU(2)_L \times U(1)$ symmetry group. The photon remains massless. Even though the former bosons have been discovered, the Higgs boson itself has eluded detection so far.

The quarks interact via the gauge bosons of the $SU(3)$ group, called gluons. There are eight gluons in total, which interact with each other due to the non-abelian nature of the $SU(3)$ group. The observed strongly interacting particles in nature are called hadrons, which can be classified into mesons (quark-antiquark states) and baryons (three quark states). In the Standard Model, quarks have an extra internal degree of freedom, namely color. Particles in nature have to consist of either a colored and an anti-colored particle, or three differently colored particles. Thus, only color singlets
### Table 1.1: Mass and charge properties of leptons.

<table>
<thead>
<tr>
<th>Lepton</th>
<th>Mass (MeV/c²)</th>
<th>Charge</th>
</tr>
</thead>
<tbody>
<tr>
<td>Electron (e)</td>
<td>0.511</td>
<td>-1</td>
</tr>
<tr>
<td>Electron Neutrino (νₑ)</td>
<td>&lt; 3·10⁻⁶</td>
<td>0</td>
</tr>
<tr>
<td>Muon (μ)</td>
<td>105.7</td>
<td>-1</td>
</tr>
<tr>
<td>Muon Neutrino (νₘ)</td>
<td>&lt; 0.19</td>
<td>0</td>
</tr>
<tr>
<td>Tau (τ)</td>
<td>1784</td>
<td>-1</td>
</tr>
<tr>
<td>Tau Neutrino (νₜ)</td>
<td>&lt; 18.2</td>
<td>0</td>
</tr>
</tbody>
</table>

### Table 1.2: Mass and charge properties of quarks. For the three lightest quarks (u, d and s) the MS masses are quoted; for the other three the pole mass is listed.

<table>
<thead>
<tr>
<th>Quark</th>
<th>Mass (GeV/c²)</th>
<th>Charge</th>
</tr>
</thead>
<tbody>
<tr>
<td>Down (d)</td>
<td>5·10⁻³</td>
<td>-8.5·10⁻³</td>
</tr>
<tr>
<td>Up (u)</td>
<td>1.5·10⁻³</td>
<td>-4.5·10⁻³</td>
</tr>
<tr>
<td>Strange (s)</td>
<td>80·10⁻³</td>
<td>-155·10⁻³</td>
</tr>
<tr>
<td>Charm (c)</td>
<td>1.5</td>
<td>1.8</td>
</tr>
<tr>
<td>Bottom (b)</td>
<td>4.6</td>
<td>5.1</td>
</tr>
<tr>
<td>Top (t)</td>
<td>169</td>
<td>179</td>
</tr>
</tbody>
</table>
exist and no physical colored states occur. The quarks interchange color through the strong interaction, with the gluons carrying the color. The strength of the strong interaction is specified by the strong coupling parameter $\alpha_s$. This parameter decreases with increasing $|Q^2|$, the absolute squared four-momentum transfer in the collision. Cross sections in QCD can be approximated by expansions in $\alpha_s$. If $\alpha_s$ is sufficiently small, which is the case for $|Q^2|$ higher than a few GeV$^2$ (so called asymptotic freedom), one can use perturbation theory to calculate the cross sections predicted by QCD reliably. For smaller $Q^2$ values, these cross sections cannot be calculated as reliably. The center of mass energy of the Tevatron, $\sqrt{s} = 1.96$ TeV, is high enough to allow precise measurements of the QCD predictions, especially for the production of bottom and charm quarks.

### 1.1 Asymptotic Freedom

In QCD the leading order quark production processes, as are shown in figure 1.1, is modified by higher order loop diagrams in which the loops consist of quarks or gluons (figure 1.2). The effect of these loops causes the bare color charge of the quark to be screened, which affects the coupling $\alpha_s$ of other quarks and gluons to the bare color charge:

$$\alpha_s(|Q^2|) = \frac{\alpha_s(m_Z^2)}{1 + (\alpha_s(m_Z^2)/12\pi)(11n - 2f) \ln(|Q^2|/m_Z^2)}$$

Here, $m_Z$ is a reference scale, in this case the Z-mass, $n$ is the number of colors, and $f$ is the number of flavors participating in the process at the given $|Q^2|$, for a certain renormalization scheme. Since in the Standard Model $n = 3$ and $f = 4$ for $b\bar{b}$ production and for the most common renormalization schemes, $11n - 2f$ is positive. Consequently, the loops have in fact an anti-screening effect on the bare color charge: with increasing $|Q^2|$, the effective coupling decreases. This is called “asymptotic freedom”. When $|Q^2|$ gets large, the quarks and gluons decouple and behave like almost free particles.

The perturbation theory on which QCD is based breaks down when higher order corrections are calculated in the ultraviolet region, corresponding to large loop momenta. Through the choice of the renormalization procedure, one can absorb the infinite parts of the corrections into the masses and coupling constants of the theory, which requires the introduction of a renormalization scale $\mu_R$. The underlying idea here is that there is a set of bare parameters that are divergent, such as the mass and the coupling constant. These bare parameters are unmeasurable. They can be renormalized such that they exactly cancel the infinities resulting from higher order corrections in any calculation. After these infinities are cancelled, the renormalized parameters remain, which obtain physical meaning and can be measured. One has to note here that the renormalization parameter $\mu_R$ is purely a mathematical tool which is needed
for the process of renormalization, and no physical consequences can result from the choice of the value of $\mu_R$. This implies that changes in cross sections due to the explicit dependance on $\mu_R$ have to be offset by changes in the renormalized parameters, such as the coupling constant and, possibly, the mass.
1.2 Heavy quark production

In the regime of asymptotic freedom, we can decompose the proton-antiproton interaction into a set of parton-parton interactions. The cross section for a certain interaction can be written as:

$$\sigma(|Q^2|) = \sum_{ij} \int dx_i dx_j \hat{\sigma}_{ij}(x_i p_A, x_j p_B, \mu_R, \mu_F, \alpha_s(\mu_R)) F_i^A(x_i, \mu_F) F_j^B(x_j, \mu_F)$$

(1.2)

where $\hat{\sigma}_{ij}$ is the parton level cross section between the partons in each hadron, $x_i$ and $x_j$ are the fractions of the proton momentum $p_A$ and anti-proton momentum $p_B$ that the partons $i$ and $j$ carry. $F_i^A$ and $F_j^B$ are the parton distribution functions (PDFs) describing the momentum distribution of the partons inside the proton and anti-proton. These structure functions contain a dependence on a factorization scale $\mu_F$, analogous to $\mu_R$, which results from the following: naively, the parton model describes the proton as containing three almost free constituents, with no interactions between them. A probe therefore scatters on a single, free and effectively massless constituent. This naive model therefore has no implicit dependence on the absolute squared four-momentum transfer $|Q^2|$ in the collision. However, the quarks inside the proton do have interactions through gluon exchange and emission. Up to a certain scale, namely the factorization scale $\mu_F$, these interactions are absorbed into the parton distributions, and therefore removed from the single parton interaction cross sections $\hat{\sigma}_{ij}$. For bottom quark production, this factorization scale is normally chosen to be a function of the mass of the $b$-quark, $m_b$, and its transverse momentum.$^1$

In QCD, the cross section of $b$-quark production in the regime of asymptotic freedom can be written as an expansion in $\alpha_s$:

$$\hat{\sigma}_{ij} = \alpha_s^2(\mu) G_{ij}^{(0)}(\hat{s}, m_b) + \alpha_s^3(\mu) G_{ij}^{(1)}(\hat{s}, m_b) + O(\alpha_s^4) \quad (\mu = \mu_F = \mu_R)$$

(1.3)

where $G_{ij}^{(0)}$ and $G_{ij}^{(1)}$ are dependent on the quark mass $m_b$ and the center of mass energy $\hat{s}$. The total cross section for $b$-quark production results from the sum of all contributing processes. The Feynman diagrams for the leading order processes are shown in figure 1.1. The experimental signature for such an event is two jets in the detector, with a transverse energy balance. Some next to leading order Feynman diagrams are shown in figures 1.2 and 1.3. The processes shown in figure 1.3 have three partons in the final state, and are either radiative corrections to the leading order processes (where we have two quarks or two gluons in the initial state), or a new class of processes in which a quark and a gluon are in the initial state.

In the following, we will study the production of $b$-jets instead of the production of $b$-quarks. The main difference between these studies is that in the case of production

$^1$The transverse momentum $p_T$ is defined as $p_T = p \sin(\theta)$, with $\theta$ the angle of the particle with respect to the beam. Similarly, the transverse energy $E_T$ is defined as $E_T = E \sin(\theta)$. 
of $b$-quarks, one is interested in the properties of the heavy quark itself, regardless of the event structure in which the $b$-quark is produced, while for a jet one is interested in the energy of the jet that contains one or more heavy quarks, and not in the fraction of the energy carried by these quarks. Calculating the properties of the quark brings with it the difficulty of the emission of collinear gluons, which at high momentum cause large logarithms to appear. These logarithms need to be resummed, and can thus be included in the fragmentation functions. The calculation of jet properties is not sensitive to the details of the analysis of large logarithms, since the jet contains both the quark and the collinear gluons. The experimental measurement of the transverse jet energy distribution therefore does not depend on the details of the fragmentation of the $b$-quark. Fragmentation describes the emissions and absorptions of gluons between the partons in the final phase of the collisions, and the splitting of gluons into quarks. In the hadronization stage the quarks and gluons, which all carry a color charge, combine to form colorless hadrons. An experimental jet is now defined as the assembly of all these hadrons, around a common axis. This is illustrated in figure 1.4.

The single particle inclusive next to leading order $b$-quark production cross section has been calculated by Nason, Dawson and Ellis [23, 24] as well as Beenakker et al. [25, 26]. Mangano and Frixione have extended this work to cover $b$-jets [27]. A theoretical calculation, providing the 4-momenta of the $b$-quark, $\bar{b}$-quark and a possible gluon, with the appropriate event weight is provided by Mangano, Nason and Ridolfi [28]. Our goal is now to use this calculation to measure the production rate for $b$-jets, that is, jets that contain one or more $b$-quarks.

Many algorithms are available that can be used to construct jets. In this analysis, we adopt the Snowmass convention, where the particles are clustered in $(\eta, \phi)$-space in a cone of radius $R$, with $R = 0.5$ throughout the analysis [29]. We only consider jets that contain a $b$ or $\bar{b}$-quark, or both; the gluon jet is disregarded if the gluon does not end up in the same jet with a heavy quark.

We calculate the cross section for $b$-jets with $|\eta_{\text{jet}}| < 0.6$, to reflect a realistic

Figure 1.3: Some diagrams showing the next to leading order corrections to heavy quark production.
1.2 Heavy quark production

Figure 1.4: Illustration of the jet definition on theoretical level and experimental level. At the theory level, we only use the heavy quark, with a possible gluon. The experimental jet takes into account the entire fragmentation and hadronization process. The out of cone showering, shown by the black lines pointing outside the jet cone, is corrected for by using the Monte Carlo simulation at a later stage.

The out of cone showering, shown by the black lines pointing outside the jet cone, is corrected for by using the Monte Carlo simulation at a later stage.

The geometrical and trigger acceptance of the DØ detector during the start of Run II. The factorization scale $\mu_F$ and renormalization scale $\mu_R$ are chosen as:

$$\mu_F^2 = \mu_R^2 = \mu_0^2 = \frac{(p_T^b)^2 + (p_T^b)^2}{2} + m_b^2 \quad (1.4)$$

and we vary these scales between $\mu_0/2$ and $2\mu_0$ to account for theoretical uncertainties. The parton distribution function adopted is CTEQ6M [30]. For the $b$-quark mass, we use $4.8 \pm 0.2$ GeV/$c^2$, with the errors reflecting the uncertainty on the mass measurement. This results in the total $b$-jet cross section as a function of $E_T^{jet}$, shown in figure 1.5. The dotted lines show the error resulting from floating $\mu_F$ and $\mu_R$, and varying the $b$-quark mass. Figure 1.6 shows the relative contributions to the full theoretical error of floating $\mu_F$ and $\mu_R$ and varying the $b$-mass. Clearly the uncertainty on the factorization and normalization scales dominates the error on the calculated cross section.

Figure 1.7 shows the theoretical cross section subdivided in the contributions of the different initial states. At high $E_T^{jet}$, the gluon splitting processes become dominant.

This is also demonstrated in figure 1.8, which shows the cross section of jets containing only one $b$-quark (or $\bar{b}$-quark) and those containing both a $b$- and a $\bar{b}$-quark. Here as well the relative contribution of $b\bar{b}$-jets increases with higher $E_T^{jet}$. 

Theoretical and experimental jet definitions on theoretical level and experimental level.
Figure 1.5: Theoretical $b$-jet production cross section, measured in $|\eta_{b-jet}| < 0.6$ and with 0.5 cone jets, with the theoretical errors resulting from floating $\mu_F$ and $\mu_R$, as well as the $b$-quark mass. The PDF set used is CTEQ6M.

Figure 1.6: Relative contribution to the error on the theoretical prediction, coming from the renormalization and factorization scales (solid line) and the uncertainty on the $b$-quark mass (dashed line). They do not add up to one due to correlations between the two.
1.2 Heavy quark production

Figure 1.7: Theoretical $b$-jet production cross section, measured in $|\eta_{b-jet}| < 0.6$ and with 0.5 cone jets, separated in the contributions of the different initial states. The PDF set used is CTEQ6M.

Figure 1.8: Theoretical $b$-jet production cross section, measured in $|\eta_{b-jet}| < 0.6$ and with 0.5 cone jets, separated in the contributions of jets with one $b$- or $\bar{b}$-quark present, or with both a $b$- and a $\bar{b}$-quark present. The PDF set used is CTEQ6M.
1.3 \( b \)-jets as a precursor to top and Higgs physics at the Tevatron

The proper understanding of the properties of \( b \)-jets is crucial for the detection and reconstruction of particles that decay into \( b \)-quarks, of which the most interesting are the top quark and the Higgs boson. Figure 1.9 shows the cross sections for \( b \), \( t \) and Higgs production at hadron colliders as a function of center of mass energy. Due to the low cross section of top and Higgs production compared to the total inelastic, non-diffractive cross section, one requires \( b \)-jets as the sensitive probes needed to sift the Higgs- and top-producing events from the background.

With the available center of mass energy of 1.96 TeV, the Tevatron currently is the only place in the world where top quark pairs can be produced directly. They decay primarily through \( t\bar{t} \rightarrow (W^+b)(W^-\bar{b}) \), resulting in two \( b \)-jets in the event. If properly detected, these \( b \)-jets can be used to remove backgrounds to the \( t\bar{t} \) signal, which gives a considerable improvement in the signal over background ratio.

The best option to detect the Higgs boson at the Tevatron is the associated production of a Higgs boson with a \( W \) or \( Z \) boson. The \( W \) and \( Z \) then decay leptonically or hadronically. The Higgs boson primarily decays to a \( b\bar{b} \) pair, since the Higgs coupling to a fermion is proportional to the squared mass of that fermion, and the \( b \)-quark is almost three times heavier than the next heavy particle (the \( \tau \) lepton). This gives rise to two \( b \)-jets in the event, in addition to the decay products of the \( W \) or \( Z \) boson.

The detection of both the top quark and the Higgs boson thus relies on proper detection and reconstruction of \( b \)-jets. The measurement of the \( b \)-jet cross section in the following chapters improves our understanding of these jets, paving the way for a better understanding of the Standard Model.
Figure 1.9: Energy dependence of interesting physics processes at hadron colliders as function of the center of mass energy. The discontinuities in the lines are caused by the change from \( \bar{p}p \) collisions to \( pp \) collisions.