Production and Decay of Charged Intermediate Vector Bosons in $p\bar{p}$ Collisions

Wisse van de Guchte
Production and Decay of Charged Intermediate Vector Bosons in $p\bar{p}$ Collisions
Stellingen behorende bij het proefschrift:

"Production and Decay of Charged Intermediate Vector Bosons in p\bar{p} Collisions"

door Wisse van de Guchte

1. De totale breedte van het $W^{\pm}$ boson kan worden afgeleid uit metingen van de partiële leptonische werkzame doorsneden voor $W^{\pm}$ en $Z^{0}$ produktie in proton-antiproton botsingen en de totale breedte van het $Z^{0}$ boson.

_Dit proefschrift, formule (2.5.2)._

2. Het oplossend vermogen van een electromagnetische calorimeter voor de positie van een deeltje is in het centrale gedeelte dikkwils gelimiteerd door de grootte van de elementaire calorimeter cel. In de voorwaartse richting leidt de verdeling van een energie depositie over meerdere cellen in het algemeen tot een betere plaats bepaling. Dit effect kan ook in het centrale gedeelte worden verkregen door een diepte segmentatie te kiezen, waarbij een toren van calorimeter cellen niet naar de bundelas wijst. Het gebruik van een zogenaamde ‘non-pointing’ toren structuur verdient daarom de voorkeur, mits nadelige gevolgen van energie verlies in cracks kunnen worden beperkt.

_J. Dorenbosch, W. van de Guchte en C. Wulz, UA1 Technical Note 87/80._

3. Uit een studie van het proces: $p\bar{p} \rightarrow W^{\pm} + 2$ jets, concluderen Mangano en Parke dat de differentiële werkzame doorsnede met betrekking tot de azimutale hoek tussen de twee jets een maximum heeft bij 180°. Dit effect heeft echter geen dynamische oorsprong, maar is slechts het gevolg van de selectie van gebeurtenissen met zeer hoge transversale energie van de jets.

_M. Mangano en S. Parke, Fermilab-Pub-89/106-T._

4. De binnen de astrologie veel gehoorde redenering, dat hemellichamen ons beïnvloeden door middel van onbekende krachten, is niet weerlegbaar. Zij heeft dientengevolge geen betrekking op de werkelijkheid.

_R. Kousbroek, “Einstein's popenhuis”, Meulenhoff Amsterdam, 1990._

5. Bij de bepaling van de koppelingen constante van de sterke wisselwerking, $\alpha_s$, met behulp van K-factoren in het proces: $p\bar{p} \rightarrow W^{\pm} + n$ jets, waarbij $n = 0$, 1 en 2, door Ansari et al. is onvoldoende rekening gehouden met de vorm van het transversaal momentum spectrum van de $W^{\pm}$ bosonen.

6. De audio data bit-rate reduktie techniek gebaseerd op het maskeer effect kan ook worden aangewend om data additie te bereiken. Het audio signaal vormt zo een “Hidden Channel”.

_W.R.Th. ten Kate, L.M. van de Kerkhof en F.F.M. Zijderveld, Rundfunktechnischen Mitteilungen Vol. 35, 10 (1991)._ 

7. Het aantal kubieke meters ‘groot’ in Nederland kan drastisch worden vermeerderd met behoud van oppervlakte door (her)invoering van de sequoia.

8. Hoewel reproduktie binnen de beeldende kunst en de muziek in eerste instantie een toenadering teweeg brengt tussen kunstenaar en publiek, leidt de daaropvolgende toename in de consumptie doorgaans tot verwijdering tussen beide.
Production and Decay of Charged Intermediate Vector Bosons in $p\bar{p}$ Collisions

PRODUCTIE EN VERVAL VAN GELADEN INTERMEDIAIRE VECTOR BOSONEN IN P$\bar{P}$ BOTSINGEN

(MET EEN SAMENVATTING IN HET NEDERLANDS)

PROEFSCHRIFT

TER VERKRIJGING VAN DE GRAAD VAN DOCTOR IN DE WISKUNDE EN NATUURWETENSCHAPPEN AAN DE RIJKSUNIVERSITEIT TE UTRECHT, OP GEZAG VAN DE RECTOR MAGNIFICUS PROF. DR. J.A. VAN GINKEL, VOLGENS BESLUIT VAN HET COLLEGE VAN DECANEN IN HET OPENBAAR TE VERDEDIGEN OP WOENSDAG 4 SEPTEMBER 1991 DES MORGENS TE 10.30 UUR

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1 Introduction
1. Introduction
For many years, electromagnetic interaction processes have been successfully described by a gauge theory, quantum electrodynamics or QED, in which elementary fermions interact by the emission or absorption of the electromagnetic ‘intermediate vector boson’ (IVB), which is the quantum representation of the electromagnetic field. Here the IVB is the photon, γ.

The gauge theory of QED has been extended in the sixties by Glashow, Salam and Weinberg\textsuperscript{1} (the GSW model), and others, to include weak interaction processes by the transfer of a triplet of ‘weak’ IVBs (the W\textsuperscript{+}, W\textsuperscript{−} and Z\textsuperscript{0} particles). A similar gauge theory, quantum chromodynamics (QCD), has been developed to describe strong interaction processes by the transfer of an octet of ‘strong’ IVBs or gluons, which transfer ‘strong’ or ‘colour’ charge. The two gauge theories form what is commonly referred to as the Standard Model.

The proton-antiproton (p\overline{p}) collider project at the Super Proton Synchrotron (SPS) at CERN has greatly contributed to the understanding of the electroweak and strong interactions between elementary particles and has provided much experimental support for the Standard Model. It was initiated in 1976 by C. Rubbia and collaborators\textsuperscript{2} after the discovery of weak neutral currents at Gargamelle at CERN\textsuperscript{3}. These currents are associated with the transfer of Z\textsuperscript{0} bosons and had been predicted by the GSW model. Weak charged current processes connected to the exchange of W\textsuperscript{±} bosons, such as for instance neutron β decays, were already known long before the model had been developed. The aim of the collider project was to create weak IVBs directly in the experiment at a time when none of the existing accelerators was capable of providing a sufficiently high center-of-mass (CM) energy for their production. The fixed target experiments at CERN and at Fermilab using 400 GeV proton beams could provide a CM energy of only 28 GeV, while the GSW theory predicted that more than 70 GeV was needed. By using a counter-rotating antiproton beam in the SPS as target, it appeared to be possible to boost the CM energy to 546 GeV! However, at that time it was far from obvious whether enough antiprotons could be created to yield the high collision rates needed for rare processes such as weak IVB production. Thanks to a technique known as stochastic beam cooling developed by S. van der Meer\textsuperscript{4} this problem could be overcome, which finally led to a successful operation of the SPS p\overline{p} collider mode (Sp\overline{p}S) under the required conditions.

Despite the pessimistic view shared by many physicists, that W\textsuperscript{±} and Z\textsuperscript{0} particles could never be seen in a ‘dirty’ environment typical for p\overline{p} collisions, the first candidates were observed in 1983 by the UA1 and UA2 collaborations\textsuperscript{5}. In the subsequent period up to 1985 (the p\overline{p} CM energy had been increased to 630 GeV) the two experiments collected several hundreds of W\textsuperscript{±}s and a few tenths of Z\textsuperscript{0}s, which allowed a more quantitative examination of Standard Model predictions concerning weak IVB production and decay. In addition to electroweak phenomena, also strong interaction processes had been observed such as the production of jets resulting from hard collisions between the proton and antiproton constituents. From these observations, QCD predictions could be tested extensively.

Although many aspects of the Standard Model were studied and became better understood, some physics topics remained unrevealed and new questions arose. For instance, there are strong reasons to believe that a sixth quark, the top quark, should exist, but it had (and has) not yet been found. Also, there are no experimental facts which force us to go beyond the Standard Model, but the fact that it contains many arbitrary assumptions and parameters leads some physicists to the belief that a more fundamental theory than the
Standard Model is required. Several candidates for such theories exist (for instance Super
Symmetry\textsuperscript{6}), and most introduce a set of new elementary particles. In the hope to cover
some of these topics and encouraged by the achievements of the running period in the early
eighties, CERN has undertaken a major upgrade of the \( p\bar{p} \) collider complex in the years 1985
to 1987 with the intention of increasing the interaction rate by an order of magnitude. After a
technical run in 1987, this goal has nearly been achieved during the end of a running period
lasting from the autumn in 1988 until the summer in 1989.

In the spring of 1987 CERN had lost its monopoly of \( p\bar{p} \) collider physics, since a
new machine had come online, the Tevatron at Fermilab (FNAL), running at a \( p\bar{p} \) CM
energy of 1.8 TeV. After an initial slow start, its performance slowly improved with time
and finally allowed the CDF experiment to collect a data sample comparable in size to the
ones gathered by UA1 and UA2. The CDF experiment is in a more favorable situation, since
cross sections of rare processes generally rise with increasing \( p\bar{p} \) CM energy.

Although much has been learned at \( p\bar{p} \) colliders, indeed the physics environment of
an \( e^+e^- \) collision is much cleaner than the one of a \( p\bar{p} \) collision. Moreover, electrons and
positrons have a pointlike nature which yields each interaction a well defined initial state,
whereas protons and antiprotons are composite objects complicating the analysis of a \( p\bar{p} \)
collision. Therefore, in many cases \( e^+e^- \) colliders are better suited to perform precision
measurements. In 1989, two new machines of this type came online, the SLC at SLAC and
LEP at CERN, both designed to measure the mass and width of the \( Z^0 \) boson with a high
accuracy. The four LEP experiments have already collected more than \( 10^5 \) \( Z^0 \)'s each and
will collect several millions in the coming years.

A second phase for the LEP collider is scheduled in which the \( e^+e^- \) CM energy is
increased from the \( Z^0 \) mass to more than twice the W mass allowing a detailed study of the
charged members of the weak IVB triplet. However, LEP II is planned to start somewhere in
1994 and until that time W bosons remain the domain of \( p\bar{p} \) colliders, because the cross
section for single W production: \( e^+e^- \rightarrow e^+W^\pm + \nu \) below the 2W threshold is negligible\textsuperscript{7}
and there is not enough phase space for the production of a W pair.

In this thesis, we will focus on the total width of the W boson, \( \Gamma_W^{\text{tot}} \). A direct mea-
surement of this quantity exists\textsuperscript{8}, but with an uncertainty of about 70%. Here, an indirect
method is presented in which \( \Gamma_W^{\text{tot}} \) is inferred from the ratio R of the W \( \rightarrow \ell\nu \) and Z \( \rightarrow \ell\ell \)
partial cross sections measured at \( p\bar{p} \) colliders and the total width of the Z boson, \( \Gamma_Z^{\text{tot}} \),
accurately measured at LEP. It exploits the complementary results of \( p\bar{p} \) and \( e^+e^- \) machines
and allows to bring down the uncertainty on \( \Gamma_W^{\text{tot}} \) to about 8%. The result will be compared
with the Standard Model prediction and its implication for the mass of the top quark will be
investigated. We will also consider the possibility to put a constraint on the masses of new
elementary particles predicted by the Super Symmetry model.

The data used for the measurement of R have been taken with the UA1 detector in the
1988 and 1989 collider runs, while for \( \Gamma_Z^{\text{tot}} \) we have used the measurement coming from
LEP. To improve the sensitivity of the result, the measurement of R has been combined with
an earlier one obtained from old UA1 data taken in the period 1983–1985 and with
measurements from the UA2 and CDF collaborations.
The outline of this thesis is as follows: in the next chapter we motivate the measurement of $\Gamma_W^{\text{tot}}$ and introduce the method employed, chapters three and four deal with the theory of weak IVB decay and production respectively, the UA1 apparatus is described in chapter five, in chapter six the simulation of $pp$ interactions in the UA1 detector is discussed, in chapter seven the analysis of the 1988 and 1989 data is presented and in chapter eight we discuss the results.

We conclude this introduction with a practical remark: for matters of convenience the speed of light is defined to be $c \equiv 1$, so that mass and momentum can always be given in units of GeV instead of GeV/$c^2$ and GeV/$c$ respectively.
2 Motivation for the Determination of $\Gamma_w^{\text{tot}}$
2. Motivation for the Determination of $\Gamma^\text{tot}_W$
2.1 Introduction

The width $\Gamma$ of weak IVBs relates to their resonance behaviour. If one considers the scattering of two particles via an intermediate resonant state, then the energy dependence of the cross section $\sigma$ for the scattering process is given by

$$\sigma(s) \propto \frac{s}{(s-M^2)^2 + M^2\Gamma^2},$$

where $s$ is the CM energy of the two particles squared, and $M$ and $\Gamma$ are the mass and width of the resonance respectively. This particular form is known as the ‘Breit-Wigner form’ and is simply the Fourier transform of an exponential time pulse, corresponding to the decay of the resonance. The resonance curve has a bell shape and is illustrated in figure 2.1.

![Breit-Wigner shape of a resonance curve.](image)

The importance of measuring the width of weak IVBs is given by the following qualitative argument: the width is actually a measure of the uncertainty in the energy, which, according to the uncertainty principle, is inversely related to time: the wider the curve, the shorter the lifetime $\tau$ of the IVB ($\tau = h/\Gamma$ where $h$ is Planck’s constant reduced). The lifetime in turn indicates the number of decay modes available to the IVB: the more particles the IVB can decay into, the briefer its lifetime—and the wider the curve—will be. Since IVBs couple to elementary fermions, a measurement of their widths might limit the number of decay modes and possibly the number of elementary fermions.

The total width of the $Z^0$ has been precisely measured at LEP with a relative uncertainty of 0.8% \(^9\). Within the context of the Standard Model (SM) its value strongly favours the existence of only three generations of fermions, although a fourth generation with heavy fermions cannot be excluded. Moreover, it seems to be consistent with a top quark (the sixth quark in the SM) which is too heavy to contribute to the $Z^0$ width.

The top quark, an essential ingredient in the SM, has been the object of intensive searches during the past years by many groups. So far, these searches did not lead to its discovery, but only allowed to constrain the range of possible values for its mass. Taking part in a global effort to pin down the top mass, we address the possibility to obtain a lower bound from a measurement of the total $W$ width. Strictly speaking, the method we use for its determination is not a direct measurement. It actually infers the quantity from the measurement of the ratio $R$ of $W$ and $Z$ partial cross sections and of $\Gamma_Z^{\text{tot}}$. 
2. Motivation for the Determination of $\Gamma_W^{\text{tot}}$

In extensions of the SM such as the minimal supersymmetric model$^6$, $\Gamma_W^{\text{tot}}$ receives significant contributions from non-standard decays. Our determination of $\Gamma_W^{\text{tot}}$ with a relative uncertainty of 8% to be discussed later in this thesis, should be sensitive to such contributions. Non-observation of an excess width with respect to the standard one, translates into bounds on the masses of new particles.

In the following, we review some indirect experimental evidence for the existence of the top quark, discuss results from direct and indirect top searches and consider constraints from $\Gamma_W^{\text{tot}}$ on supersymmetric W decays. We then give an outline of the method to determine $\Gamma_W^{\text{tot}}$ and conclude this chapter with a review of related measurements.

2.2 Experimental Indications for the Existence of the Top Quark

2.2.1 Introduction

The minimal SM, based on the gauge group$^{10}$ SU(3)$_c$×SU(2)$_L$×U(1)$_Y$ with three generations, is in very good agreement with the experimental data available at present. Ingredients still missing, however, are the top quark and the Higgs boson. Here, we restrict our attention to the former.

In the SM, the left-handed top quark is the upper member of a third generation weak isospin doublet: $(\bar{t}^c, t^c)_L$, in which its partner is the bottom quark. This quark doublet is accompanied by a third generation lepton doublet, made up of the tau and the tau-neutrino: $(\bar{\tau}^c, \tau^c)_L$. The right-handed states are organized in weak isospin singlets: $t_R$, $b_R$ and $\tau_R$, with an exception for $\nu_\tau$ of which no right-handed version exists. Experimental evidence for a top quark comes from the observation of the $t$ resonance$^{11}$, while direct production of the $t$-quark has been observed in $e^+e^-$ collisions$^{12}$. So far, the evidence for a $t$ is only indirect and comes from the observation of leptonic $t \to \tau$ decays. However, it is likely that the $\nu_\tau$ occupies the ‘up’-state of a weak isospin doublet and there are strong indications that indeed it is a sequential neutrino, that is, distinct from the first and second generation neutrinos: $\nu_e$ and $\nu_\mu$$^{13}$.

A purely theoretical argument for the existence of a t-quark originates from the requirement that the SM should be a renormalizable field theory, that is, one in which all observables are calculable in terms of a finite number of measured parameters. Renormalizability can be upset by the occurrence of ‘anomalies’. In electroweak theories such as the GSW model, anomalies arise when gauge bosons couple differently to left- and right-handed fermion states (which is indeed the case for $Z^0$ bosons, while W bosons do not couple to all to right-handed states). These different couplings can lead to infinities in, for instance, the calculation of the interaction of three gauge bosons via a closed fermion loop, which cannot be removed. The specific theory can be made anomaly-free only, if the sum of the separate fermion contributions to anomalies exactly cancels. For the GSW theory, this leads to the requirement that the sum over the electric charges of all particles in right-handed doublets minus the one for left-handed doublets must be zero$^{14}$. Since right-handed fermion doublets do not exist, this means that the sum over the electric charges of all particles in left-handed doublets must vanish. This can be achieved in each individual fermion family, if a lepton doublet is accompanied by a quark doublet where the quarks come in three different colours.
Thus, the observation of a third generation lepton doublet motivates the existence of a third generation quark doublet.

We consider two types of measurements that provide experimental support for top. The first reveals the $I_3 = -1/2$ nature of the b-quark, where $I_3$ is the third component of weak isospin. This value suggests that the b-quark should have a partner with $I_3 = +1/2$. The second measurement concerns mixing of $B^0$ and $\bar{B}^0$ mesons. The invariant amplitude which accounts for this phenomenon is proportional to the masses squared of quarks which occur in the intermediate state. The observed amount of mixing requires a quark mass heavier than the ones known so far.

### 2.2.2 Bottom as the Down State of a Weak Isospin Doublet

The isospin of the b-quark can be determined from the forward-backward asymmetry, $A_{FB}$, in the angular distribution of b’s produced in $e^+e^-$ annihilation. It is defined as

$$A_{FB} = \frac{\int_0^1 dz (d\sigma/dz) - \int_{-1}^0 dz (d\sigma/dz)}{\int_{-1}^1 dz (d\sigma/dz)}, \quad (2.2.1)$$

where $z = \cos\theta$ ($\theta$ is the polar angle of the b-quark with respect to the beam) and $d\sigma/dz$ is the differential cross section for the process: $e^+e^- \to b\bar{b}$, which can easily be calculated and is given by

$$\frac{d\sigma}{dz} = 3 \frac{2\pi \alpha^2}{4s} [R(1+z^2) + Bz],$$

$$R = Q_b^2 - 2Q_bv_c\bar{v}_bRe(\chi) + (a_c^2 + v_c^2)(a_b^2 + v_b^2)\chi^2,$$

$$B = 4a_c a_b[-Q_bRe(\chi) + 2v_c\bar{v}_b|\chi|^2], \quad (2.2.2)$$

$$\chi = \frac{1}{4\sin^2\theta_W\cos^2\theta_W} \frac{s}{s - M_Z^2 + iM_Z\Gamma_Z},$$

where $\alpha$ is the fine structure constant, $s$ the electron-positron CM energy squared, $Q_b$ the b-quark charge and $\sin^2\theta_W$ the weak mixing parameter. The quantity $R$ is the total cross section for $b\bar{b}$ production normalized to the purely electromagnetic cross section for the process: $e^+e^- \to \mu^+\mu^-$. It contains three terms, of which the first and last correspond to pure $\gamma$ and $Z^0$ exchange respectively, while the middle one is associated with $\gamma - Z^0$ interference. The $\chi$ function is a Breit-Wigner function and describes the resonance behaviour of the $Z^0$. The asymmetry is caused by the weak interaction and is contained in the B-term. Both R and B depend on the vector and axial-vector couplings of the electron and b-quark to the $Z^0$

$$a_T = I_3^T,$$

$$v_T = I_3^T - 2Q_b\sin^2\theta_W. \quad (2.2.3)$$
2. Motivation for the Determination of $\Gamma_W^{\text{tot}}$

![Graph showing forward-backward asymmetry for $e^+e^- \to b\bar{b}$](image)

Figure 2.2: Forward-backward asymmetry for $e^+e^- \to b\bar{b}$. The data points clearly favour the value $I_3 = -1/2$ for the b-quark.

After substitution of (2.2.2) in (2.2.1) we find for the forward-backward asymmetry of b-production in $e^+e^-$ annihilation

$$A_{FB} = \frac{3B}{8R},$$

which via the relations in (2.2.3) is a direct function of the weak isospin of the b-quark: $I_3^b$. Figure 2.2 shows the SM prediction for $A_{FB}$ versus $\sqrt{s}$ for three different values of $I_3^b$: -1/2, 0 and +1/2 (for: $M_Z = 91.2$ GeV, $\Gamma_Z = 2.5$ GeV and $\sin^2\theta_W = 0.230$). The L3 measurement at LEP excludes $I_3^b = 0$, but cannot distinguish between the other two values. This is clear from the definition of B in (2.2.2): At LEP, the first term can be neglected, because pure $Z^0$ exchange dominates $\gamma - Z^0$ interference. Therefore, B is proportional to $a_bV_b$ and hence almost to $(I_3^b)^2$, so that the sign drops out. However, at intermediate and low energy where the interference term is more important, the data points clearly favour the value: $I_3^b = -1/2$.

2.2.3 Mixing in the $B^0\bar{B}^0$ System

The production of a $c\bar{c}$ or $b\bar{b}$ pair in $p\bar{p}$ or $e^+e^-$ collisions can give rise to opposite-sign lepton pairs if both heavy quarks decay semileptonically ($c \to s \ell^+\nu$ and $b \to c \ell^-\nu$). In $b\bar{b}$ events, also like-sign lepton pairs can be produced. This proceeds via cascade decays, in which c-quarks from semileptonic b-decays themselves also decay semileptonically.
2.2 Experimental Indications for the Existence of the Top Quark

An analysis of heavy flavour data by the UA1 collaboration\(^\text{17}\) revealed an excess of like-sign dimuon events over opposite-sign dimuon events which could not be explained in the conventional way. It was attributed to ‘mixing’ or ‘oscillation’ of \(B^0\) and \(\bar{B}^0\) states. If after production of a \(B^0\bar{B}^0\) pair, one of the mesons would decay semileptonically while the other would transform first into its antiparticle before decaying semileptonically (such a decay is often called an ‘oscillated decay’ as opposed to a ‘normal’ decay), the resulting charged leptons would have equal sign. The measurement was not sensitive to the type of \(B^0\) meson: \(B^0_{(s)}(bd)\) or \(B^0_{s}(\bar{b}s)\). About simultaneously, the ARGUS collaboration\(^\text{18}\) observed the same phenomenon in \(e^+e^-\) collisions and could explicitly ascribe the mixing to \(B^0_d\) mesons. They quote the following probability for an oscillated \(B^0_d\) decay:

\[
\chi_d = P(B^0_d \rightarrow \bar{B}^0_d) = 0.17 \pm 0.05 \tag{2.2.5}
\]

The interactions responsible for these oscillations are second order weak interactions of the type shown in figure 2.3. If we denote the quarks in the intermediate state by \(Q\) then the amplitude is proportional to the quark mass squared and to products of Kobayashi-Maskawa (KM) matrix elements: \(m^2_QV_{Qb}V_{Qd}^*\) for \(Q = u, c\) or \(t\). Since the products of KM matrix elements for \(u\)-, \(c\)- or \(t\)-quarks are all of the same order of magnitude, the dominating factor is the quark mass. Neglecting the contributions from \(u\)- and \(c\)-quarks, which are less than 1\%, the full expression for the invariant amplitude describing a \(B^0_d \rightarrow \bar{B}^0_d\) transition is given by\(^\text{19}\)

\[
\mathcal{M} = \frac{G_F^2}{12\pi^2} B_{B^0_d}^0 m_{B_d} m_t |V_{tb}|^2 |V_{td}|^2 \eta_{QCD} \frac{A(z)}{z} \tag{2.2.6}
\]

In this expression, \(G_F\) is the Fermi coupling constant, \(B\) is a confinement factor (\(B = 1\)), \(f_{B_d}\) is the \(B^0_d\) decay constant which is of the order of 0.1 GeV, \(m_{B_d} = 5.28\) GeV is the \(B^0_d\) mass and \(\eta_{QCD}\) is a QCD correction factor which varies slowly from 0.80 for \(m_t = 40\) GeV to 0.85 for \(m_t = 200\) GeV. Finally, \(A(z)/z\) is related to higher order terms in \(z \equiv m_t^2/M_W^2\). It decreases from 1 to 1/4 as \(m_t\) goes to 0 to \(\infty\). The dominating uncertainties in relation (2.2.6) (apart from \(m_t\)) come from the KM matrix element \(V_{td}\) and the \(B^0_d\) decay constant. The oscillated decay probability is directly related to the invariant amplitude by

\[
\chi_d = \frac{1}{2} \frac{(2\mathcal{M} \tau_B)^2}{1 + (2\mathcal{M} \tau_B)^2}, \tag{2.2.7}
\]
2. Motivation for the Determination of $\Gamma^\text{tot}_W$

![Graph showing the top mass dependence of $\chi_d$ with different input values for $(\eta B_{td}, V_{td})$: a) (0.14 GeV, 0.017), b) (0.18 GeV, 0.023) and c) (0.10 GeV, 0.001), and the ARGUS result.](image)

Figure 2.4: The top mass dependence of $\chi_d$ with different input values for $(\eta B_{td}, V_{td})$: a) (0.14 GeV, 0.017), b) (0.18 GeV, 0.023) and c) (0.10 GeV, 0.001), and the ARGUS result.

where $\tau_B$ is the B meson lifetime ($0.12 \times 10^{-11}$ sec). Figure 2.4 shows the $m_t$-dependence of $\chi_d$ for three cases, in which the largest unknowns in relation (2.2.6) have been chosen such that they yield two extreme and a central value for $\chi_d$. A comparison with the ARGUS result shows that top masses smaller than about 40 GeV are unlikely.

2.3 Search for the Top Quark

2.3.1 Introduction

After having given some experimental indications that support the existence of a top quark, we now consider explicit top searches and efforts to determine the top mass, $m_t$. Direct searches have been performed both at $\bar{p}p$ and $e^+e^-$ colliders. At present, however, the search at $e^+e^-$ machines is limited to a mass range for top that cannot exceed a kinematic limit of $M_Z/2$. Therefore, we only consider results from hadron colliders that go beyond that limit. Since one usually tries to detect the t-quark by tagging the t-jet with a lepton, the lower limits on $m_t$ resulting from these searches always depend on the assumption that the semileptonic branching ratio $Br(t \rightarrow b/\ell\nu)$ equals 1/9.

The top mass also enters SM predictions for electroweak parameters through radiative corrections. In some cases, the $m_t$-dependence is rather strong and, if precision measurements of these parameters are available, can be used to probe $m_t$ indirectly even if the t-quark is too heavy to be produced directly in $\bar{p}p$ or $e^+e^-$ collisions.
Finally, if the top is light enough, the decays: $Z^0 \rightarrow t\bar{t}$ and $W \rightarrow t\bar{b}$ should give measurable contributions to the total widths of the weak IVBs. The absence of such contributions implies that the t-quark must be heavy. Contrary to direct searches, the limits on $m_t$ derived here do not depend on any assumption. Even though $\Gamma_{t\bar{t}}^{\text{tot}}$ has been measured much more precisely than $\Gamma_{W}^{\text{tot}}$, an explicit motivation for the latter is that the decay $W \rightarrow t\bar{b}$ allows to cover a larger range in $m_t$ than the decay $Z^0 \rightarrow t\bar{t}$ owing to the small b-mass.

2.3.2 Direct Searches at Hadron Colliders

The production of t-quarks in pp collisions can proceed along two ways:

(i) Via the strong interaction: $p\bar{p} \rightarrow t\bar{t} + X$, mainly by quark-antiquark annihilation and by gluon-gluon fusion. A next-to-leading order QCD calculation for this process exists\(^{20}\), which allows to give predictions for the $t\bar{t}$ production cross section with an uncertainty of about 30\%\(^{21}\) (mainly due to uncertainties on the structure functions, QCD scale and $\alpha_S$).

(ii) Via an intermediate W: $p\bar{p} \rightarrow W + X$ which subsequently decays: $W^+ \rightarrow t\bar{b}$ or $W^- \rightarrow \bar{t}b$. The cross section for this process is much better known than the previous one, because it can be related to a partial leptonic cross section by

$$\sigma(p\bar{p} \rightarrow W^+X \rightarrow t\bar{b}X) = 3\sigma(p\bar{p} \rightarrow W^+X \rightarrow \ell^+\nu X)\text{PS}(m_t)\eta_{\text{QCD}}(m_t),\quad (2.3.1)$$

where the factor 3 is a colour factor, PS($m_t$) is a known phase space suppression factor and $\eta_{\text{QCD}}(m_t)$ contains QCD corrections which are of the order of a few percent (both PS and $\eta$ will be considered in the next chapter). The cross section times leptonic branching ratio can be replaced by measurements of this quantity by the UA2\(^{22}\) and CDF\(^{23}\) collaborations at $p\bar{p}$ CM energies ($\sqrt{s}$) of 0.63 TeV and 1.8 TeV respectively. The former obtained a relative uncertainty of 6\%, the latter about 17\%.
2. Motivation for the Determination of $\Gamma_w^{\text{tot}}$

The relative importance of the two production mechanisms depends on $m_t$ and on $\sqrt{s}$. This is illustrated in figure 2.5, which shows the $m_t$ dependence of the strong and weak production cross sections at CERN and at FNAL energies. At $\sqrt{s} = 0.63$ TeV, one can distinguish three distinct dynamical regimes. For $m_t \lesssim 45$ GeV most t-quarks are produced strongly, while for $45$ GeV $\lesssim m_t \lesssim M_W$ the weak production mechanism is dominant. Beyond $M_W$ weak production is no longer kinematically allowed, so that only $t\bar{t}$ production remains. At $\sqrt{s} = 1.8$ TeV, most t-quarks are produced strongly for all values of $m_t$.

In the SM with three generations the KM matrix element $V_{tb} \approx 99.9\%$, so that all t-quarks decay through virtual W exchange (or through real on-shell W’s for $m_t \gtrsim M_W$) into: $t \to Wb \to q\bar{q}b$, or: $t \to Wb \to \ell\nu b$. Therefore, the final state will contain at least four jets for the first process or a charged lepton and at least two jets for the second process. At $p\bar{p}$ colliders, it is not possible to detect the four-jet final state, because of the overwhelming background from other QCD processes. One is thus limited to the semileptonic decay mode for which a branching ratio of $1/9$ is assumed.

Recently, top searches have been reported by the UA1$^{24}$, UA2$^{25}$ and CDF$^{26}$ collaborations. The UA1 and UA2 experiments have a muon and electron detection capability respectively, while CDF can detect both muons and electrons.

The experiments at the SpS are sensitive to top masses up to $\sim 65$ GeV. In this region, where weak production dominates, the final state consists of $\ell\nu b$ balanced in the transverse plane by $b$; the charged lepton is expected to be isolated owing to the large top mass; the $\ell\nu$ mass reflects the $W$ mass. Two main types of background can be expected: (i) strongly produced $b\bar{b}$ pairs, where one the quarks decays semileptonically, (ii) $W + n$ jets, where the $W$ decays directly into a $\ell\nu$ pair. UA1 has executed its search in $(\mu + \geq 2\text{jet})$ and $(\mu\mu + \geq 1\text{jet})$ events. Among the selection criteria is a cut on the reconstructed $\mu\nu$ mass which restricts it to low values. This eliminates the second type of background. The analysis strongly relies on the muon isolation, but also employs other variables. Figure 2.6 shows the 90% C.L. and 95% C.L. limits on the top production cross section that result from their search, and the theoretical prediction. Top masses below 60 GeV are excluded at the 95% C.L. The search executed by UA2 in $(e + \geq 1\text{jet})$ events is complementary to the one of UA1, because their selection criteria include cuts on the reconstructed transverse momenta of the electron and the neutrino which effectively restricts ev masses to large values. Therefore, the main background they have to consider comes from $W + \text{jet}$ production. The events passing their selection can be fully explained in terms of this background, from which they conclude that top masses below 69 GeV are excluded at the 95% C.L.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>Limit at 95% C.L.</th>
</tr>
</thead>
<tbody>
<tr>
<td>UA1$^{24}$</td>
<td>$m_t &gt; 60$ GeV</td>
</tr>
<tr>
<td>UA2$^{25}$</td>
<td>$m_t &gt; 69$ GeV</td>
</tr>
<tr>
<td>CDF$^{26}$</td>
<td>$m_t &gt; 89$ GeV</td>
</tr>
</tbody>
</table>
2.3 Search for the Top Quark

![Graph showing the limits on the t-quark production cross section](image)

Figure 2.6: UA1 limits on the t-quark production cross section $[\sigma(\bar{t}) + \sigma(t)]$ as a function of its mass, and the theoretical prediction.

The CDF collaboration, searching for top at $\sqrt{s} = 1.8$ TeV, covers a much larger mass range than the previous two experiments, because the t-quark production cross section strongly rises with $\sqrt{s}$ as can be seen in figure 2.5. With the present luminosity (4.4 pb$^{-1}$) they are sensitive to top masses almost up to $M_Z$. The final state to be considered here is given by the process: \( p\bar{p} \rightarrow t\bar{t}X \rightarrow W^+W^-bbX \). Again, detection of the multi-jet final state is excluded, because of the large QCD background. They have searched for top in: (e-e), (e-\mu), (\mu-\mu), (e\nu+jets) and (\mu\nu+jets) events and obtain a lower bound on $m_t$ of 89 GeV at 95% C.L. using all measurements and the expected $m_t$-dependence of the cross section.

2.3.3 Constraints on the Top Mass from Global Analyses of Electroweak Data

In the minimal SM, the main unknowns among the electroweak parameters are $m_t$ and $m_h$ (Higgs boson mass). Our ignorance of $m_t$ is a serious limitation for precise tests of electroweak theory, because radiative corrections depend strongly on $m_t$ (the sensitivity of these corrections to $m_h$ is much smaller). However, one can invert the statement and exploit this dependence to get some information about $m_t$. In fact, the requirement that the electroweak theory be consistent with the data, restricts the allowed range of $m_t$ significantly.

A quantity often considered in global analyses of electroweak data is the weak mixing parameter, $\sin^2\theta_W$. Its value has been determined experimentally in different types of experiments such as measurements of the weak IVB masses, the ratio $R_V$ of neutral- to charged-current deep-inelastic neutrino scattering, parity violation in atomic caesium, etc. The strongest constraint on $m_t$ comes from the combination of results on the IVB masses and on $R_V$.

At present, the cleanest way to obtain $\sin^2\theta_W$ is from $M_Z$. There is no theoretical error associated with the measurement of $M_Z$ in e$^+e^-$ annihilation, and the only ambiguity in the derivation of $\sin^2\theta_W$ from $M_Z$ is our ignorance of $m_t$ and $m_h$. The two are related by
2. Motivation for the Determination of $I_{W}^{\text{tot}}$

\[
\sin^2 \theta_W \cos^2 \theta_W = \frac{\pi \alpha}{\sqrt{2} G_F M_Z^2 (1 - \Delta r)}.
\]

Both $\alpha$ and $G_F$ are known with 'infinite' precision ($\alpha = 1/137.0359895(61)$ and $G_F = 1.166389(22) \times 10^{-5} \text{GeV}^{-2}$) and also $M_Z$ is now accurately known from LEP$^9$: $M_Z = 91.172 \pm 0.031$ GeV. The quantity $\Delta r$ represents the radiative corrections and hides the dependence on $m_t$ and $m_h$. Its functional form depends on the renormalization scheme employed to calculate the corrections. An overview of the behaviour of $\Delta r(m_t)$ in different schemes can be found, for instance, in reference 27. A common definition is$^{28}$

\[
\sin^2 \theta_W = 1 - \frac{M_W^2}{M_Z^2}.
\]

In this scheme, the corrections depend quadratically on $m_t$ and logarithmically on $m_h$. The $m_t$-dependence of $\sin^2 \theta_W$ is illustrated in figure 2.7$^{29}$ by the curve labelled VBM for $m_h = M_Z$ (the band reflects the experimental uncertainty on $M_Z$ at a time when LEP results were not yet available). In the same figure is indicated the value of $\sin^2 \theta_W$ as obtained from direct measurements of the ratio $M_W/M_Z$ which is independent of $m_t$ by definition (the value shown in the figure is based on UA2 and CDF results presented at a 1989 conference and has slightly changed). This mass ratio can be measured cleanly at $p\bar{p}$ colliders, because the energy scale uncertainty associated with mass measurements drops out in the ratio. Note that an upper limit of $\sim 250$ GeV on $m_t$ is already implied by the IVB masses alone.

![Figure 2.7: $\sin^2 \theta_W$ obtained from $M_Z$, $M_W/M_Z$, and $R_V$
$= \sigma_{WNC}^{CC}/\sigma_{VNC}$ as a function of $m_t$ (from reference 29).](image)
2.3 Search for the Top Quark

A value of $\sin^2 \theta_W$ can also be obtained from measurements of the ratio of neutral- to charged-current cross sections, $R_V \equiv \sigma_{\text{NN}}^{\text{NC}} / \sigma_{\text{NN}}^{\text{CC}}$. A zeroth order approximation of this quantity is given by

$$R_V = g_L^2 + g_R^2 r$$

where the left- and right-handed couplings to the u- and d-quarks in the nucleons can be approximated by

$$g_L^2 = \frac{1}{2} - \sin^2 \theta_W + \frac{5}{9} \sin^4 \theta_W$$
$$g_R^2 = \frac{5}{9} \sin^4 \theta_W$$

and $r \equiv \sigma_{\text{NN}}^{\text{CC}} / \sigma_{\text{NN}}^{\text{NC}}$ is the ratio of $V$ and $V$ charged-current cross sections, which can be measured directly. The ratio $R_V$ has been measured to a 1% accuracy by the CHARM$^{31}$ and CDHS$^{32}$ collaborations. However, the extraction of $\sin^2 \theta_W$ from these measurements is affected by rather large theoretical uncertainties associated with the validity of the parton model used to describe the nucleus, the charm mass (which induces a threshold effect in the production of charmed hadrons via charged currents), etc. The value of $\sin^2 \theta_W$ derived from $R_V$ is nearly independent of $m_t$ as can be seen from the curve in figure 2.7 labelled by $\nu-q$ (the band reflects both the experimental and theoretical uncertainty).

It is clear from figure 2.7, that requiring consistency between the results obtained with these two methods leads to an upper and lower bound on $m_t$. The authors of reference 29 have combined these results (including a few more, but less constraining ones) in a global fit and obtained the following value for the top mass

$$m_t = 132^{+31}_{-37} \text{ GeV}$$

for $m_h = M_Z$. They also considered the effect of different Higgs masses and found a shift in the preferred value for $m_t$ of +5 GeV and -10 GeV for $m_h = 0.1M_Z$ and $m_h = 10M_Z$ respectively. More recently, similar attempts$^{33}$ have been made to determine the preferred range of top masses using the latest measurements of $M_Z$. However, they do not significantly improve the result quoted in (2.3.6). This can be well understood from figure 2.7 which shows that the limiting factor is $R_V$ rather than $M_Z$.

The lower bound on $m_t$ in (2.3.6) can be avoided if new heavy particles contribute to vacuum polarization diagrams, while the upper bound remains valid$^{27}$.

2.3.4 Constraints on the Top Mass from $\Gamma_{Z}^{10t}$ and $\Gamma_{W}^{10t}$

Fermions contribute to the $Z^0$ and $W$ widths via the decays

$$Z^0 \rightarrow f\bar{f}$$
$$W \rightarrow f\bar{f}$$

with partial width indicated by $\Gamma_{Z}^{ff}$ and $\Gamma_{W}^{ff}$ respectively. Their total widths are related to the number of elementary fermions for which the decays in (2.3.7) are kinematically allowed.
In the SM with three generations, the $Z^0$ total width is given by

$$\Gamma_{Z}^{\text{tot}} = \sum_{\ell} \Gamma_{Z}^{\ell\ell} + \sum_{v} \Gamma_{Z}^{vv} + \sum_{q} \Gamma_{Z}^{qq} + \Gamma_{Z}^{tt}(m_t)$$

(2.3.8)

where $\ell \equiv e, \mu, \tau$, $v = \nu_e, \nu_\mu, \nu_\tau$ and $q = u, d, c, s, b$, and the W total width by

$$\Gamma_{W}^{\text{tot}} = \sum_{\ell} \Gamma_{W}^{\ell\ell} + \sum_{q,q'} \Gamma_{W}^{qq'} + \Gamma_{W}^{tb}(m_t)$$

(2.3.9)

with $q = u, c$, $q' = d, s, b$ and similar assignments for $\ell$ and $v$. Explicit expressions for the partial widths will be given in the next chapter. Figure 2.8 shows the effect of the top mass on the partial widths $\Gamma_{Z}^{\ell\ell}$ and $\Gamma_{W}^{tb}$. The contribution of top to the W width is much larger than to the Z width for small $m_t$. For $m_t = M_Z/2$, $\Gamma_{Z}^{tt}$ has been reduced to zero while $\Gamma_{W}^{tb}$ still equals about 350 MeV. The latter becomes negligible at $m_t = M_W - m_b = 75$ GeV.

**Table 2.2:** $\Gamma_{Z}^{\text{tot}}$ as measured by the four LEP experiments\(^9\) and the SM prediction for three generations without a contribution from top.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>$\Gamma_{Z}^{\text{tot}}$ (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>ALEPH</td>
<td>2498 ± 31</td>
</tr>
<tr>
<td>DELPHI</td>
<td>2495 ± 40</td>
</tr>
<tr>
<td>L3</td>
<td>2502 ± 33</td>
</tr>
<tr>
<td>OPAL</td>
<td>2497 ± 34</td>
</tr>
<tr>
<td>Average</td>
<td>2498 ± 20</td>
</tr>
<tr>
<td>SM prediction</td>
<td>2492 ± 20</td>
</tr>
</tbody>
</table>
2.4 Supersymmetric Contributions to the W Decay Width

The total width of the $Z^0$ has been measured with an uncertainty smaller than 1% by the four experiments at LEP, for which the results$^9$ have been indicated in table 2.2. The error on the average is due to a 17 MeV statistical error and a 10 MeV systematic uncertainty common to the four experiments. Also indicated is the SM prediction assuming no contribution from a t-quark. It is affected by our ignorance of $m_t$ and $m_h$ through radiative corrections and by the uncertainty on $\alpha_S$. Clearly, there is no room left for a contribution from top. This implies that $m_t$ must be larger than $M_Z/2 \approx 45$ GeV. Thus the $m_t$-range left to be explored by measurements of the W width extends from 45 GeV to 75 GeV.

2.4 Supersymmetric Contributions to the W Decay Width

As mentioned earlier, the SM is in good agreement with the available data. However, it is widely believed to be incomplete. It contains many arbitrary assumptions and parameters. For instance, why are left-handed fermions in SU(2) doublets and right-handed ones in SU(2) singlets, why three colours? Etc. Another problem arises if one calculates the radiative corrections to the mass of the Higgs boson of the SM. These corrections are divergent, that is, they are of the form $\delta m^2 \approx \Lambda^2$ where $\Lambda$ is the scale beyond which the theory no longer applies$^{34}$. This scale is believed to be an effective cutoff scale, beyond which new physics can be expected. A detailed discussion of problems associated with the SM, which are subtle and of a theoretical nature, can be found in reference 34.

The incompleteness of the SM has motivated physicists to construct new theories (of which the SM can usually be considered as the effective theory at low energies). These new models introduce new elementary particles. One way to search for these particles is to look for unusual decay modes of the weak IVBs. The advantage of employing their decay widths over direct search methods is that one does not need to consider the topological details of specific final states. However, a clear disadvantage is that the decays must be kinematically allowed thus limiting the accessible range of new particle masses. As was shown in table 2.2, the measured value of $\Gamma_W$ is very close to the SM prediction (as are the partial hadronic and leptonic widths$^9$) leaving little room for additional $Z^0$ decays. We are therefore interested in decays which increase $\Gamma_W$ without increasing $\Gamma_Z$. This is not easy, since in most cases non-standard physics tends to contribute more to $\Gamma_Z$ than to $\Gamma_W$ $^{35}$, the classical example being additional light neutrinos. A possible exception might come from supersymmetric W decays into charginos and neutralinos$^{34}$.

In the minimal supersymmetric model (MSSM), there are two charginos, $\tilde{\chi}^\pm$ and $\tilde{\chi}'^\pm$ in increasing order of mass, and four neutralinos, $\tilde{\chi}_i$ ($i = 1,\ldots,4$) in increasing order of mass. These mass eigenstates are mixtures of weak interaction eigenstates formed by the wino ($\tilde{W}^\pm$), zino ($\tilde{Z}^0$) and photino ($\tilde{\gamma}$) (collectively called gauginos) and by the higgsinos: $\tilde{H}_{1,2}^\pm$ and $\tilde{H}_{1,2}^0$. The gauginos and higgsinos are the superpartners of gauge bosons and higgs bosons in the SM respectively.
2. Motivation for the Determination of $\Gamma_{\tilde{W}}^{\text{tot}}$

Figure 2.9: Supersymmetric contributions ($\Delta \Gamma$) to the W and Z decay widths versus the wino mass ($M_{\tilde{W}}$).

Both the W and Z$^0$ can decay into pairs of charginos and neutralinos. The decays of interest are

$$W^\pm \rightarrow \tilde{\chi}^\pm \tilde{\chi}_1, \quad W^\pm \rightarrow \tilde{\chi}^\pm \tilde{\chi}_2, \quad Z^0 \rightarrow \tilde{\chi}_1 \tilde{\chi}_2, \quad Z^0 \rightarrow \tilde{\chi}^+ \tilde{\chi}^-.$$  \hspace{1cm} (2.4.1)

The associated decay widths are dependent on the elements of the mass matrices describing the mixing of the weak interaction eigenstates. Explicit expressions are available\textsuperscript{36}, and will be given in the next chapter. To illustrate the fact that supersymmetric contributions to the widths can be large, we quote the result of an old calculation\textsuperscript{37} where the authors made a specific choice for the mass matrix parameters in which the coupling for the lightest chargino is W-like and for the two lightest neutralinos $\gamma$-like and Z$^0$-like (it is customary, to denote such particles by: $\tilde{\omega}^\pm$, $\tilde{\gamma}$ and $\tilde{Z}$ respectively). They considered the following decays: $W^\pm \rightarrow \tilde{\omega}^\pm \tilde{\gamma}$, $W^\pm \rightarrow \tilde{\omega}^\pm \tilde{Z}$ and $Z^0 \rightarrow \tilde{\omega}^+ \tilde{\omega}^-$ (shortly denoted by: $X \rightarrow ab$) and obtained the following general expression for the decay widths

$$\Gamma_X^{ab} = \frac{1}{3} M_X \alpha \lambda^{1/2}(1,x_a,x_b) \times \{ (g_v^2 + g_a^2)[1 - \frac{1}{2} x_a - \frac{1}{2} x_b - \frac{1}{2} (x_a - x_b)^2] + 3(g_v^2 - g_a^2)(x_a x_b)^{1/2} \},$$  \hspace{1cm} (2.4.2)

where $x_a = m_a^2/M_X^2$, $x_b = m_b^2/M_X^2$, $\alpha = 1/128$,

$$\lambda(1,x_a,x_b) = 1 + x_a^2 + x_b^2 - 2x_a - 2x_b - 2x_a x_b,$$  \hspace{1cm} (2.4.3)

and the vector and axial-vector coupling, $g_v$ and $g_a$, are given in table 1 of the same reference. Figure 2.9 shows these widths versus the $\tilde{\omega}^\pm$ mass where the $\tilde{\gamma}$-mass is zero and the $\tilde{Z}$-mass depends on the $\tilde{\omega}^\pm$ mass (the other input parameters are: $M_W = 80$ GeV, $M_Z = 91$ GeV and $\sin^2 \theta_W = 0.230$). The figure illustrates the possibility to obtain stronger limits on the $\tilde{\omega}^\pm$ mass from the decay $W^\pm \rightarrow \tilde{\omega}^\pm \tilde{\gamma}$, than from the decay $Z^0 \rightarrow \tilde{\omega}^+ \tilde{\omega}^-$. Unfortunately, the required sensitivity is beyond the one attained in this thesis.
Since a priori none of the couplings is more likely than others, it is better to express the widths directly in terms of the parameters which define the mass matrices in the chargino and neutralino sectors (next chapter). By comparing the measured widths with the predicted ones, one can then exclude certain regions in this parameter space. This has direct implications for the chargino and neutralino masses, because each point in parameter space corresponds to a set of values for these masses, and allows to set limits in a more systematic way. Finally, it should be remarked that the possibility for certain decays to contribute to $\Gamma_W^{\text{tot}}$ without having related decays contributing to $\Gamma_Z^{\text{tot}}$ is, of course, due to the many degrees of freedom of the MSSM.

2.5 The Method

At $p\bar{p}$ colliders, $W$ and $Z$ bosons are detected through their leptonic decay mode, that is, by the decays: $W \to \ell^+\nu$ and $Z^0 \to \ell^+\ell^-$, where $\ell = e, \mu$. The partial production cross sections are related to the $W$ and $Z$ total widths through the branching ratios for these particular decay modes, by

$$
\sigma(p\bar{p} \to WX \to \ell^+\nu X) = \sigma(p\bar{p} \to WX) \times [\Gamma_W^{\ell\nu}/\Gamma_W^{\text{tot}}],
$$

$$
\sigma(p\bar{p} \to Z^0X \to \ell^+\ell^-X) = \sigma(p\bar{p} \to Z^0X) \times [\Gamma_Z^{\ell\ell}/\Gamma_Z^{\text{tot}}].
$$

(2.5.1)

In principle, the total widths can be inferred from these relations using measurements of the partial leptonic cross sections and SM predictions for the total cross sections and the partial leptonic widths. In practice, however, the uncertainties involved are rather large. The experimental uncertainty on the partial cross sections is limited by the uncertainty on the luminosity (5% for UA2, and even larger for UA1 and CDF), while the theoretical uncertainty on the total cross sections is of the order of 8% (see chapter four).

In 1983, it was proposed\(^{38}\) to consider the ratio of the two relations in (2.5.1), because most of the theoretical and experimental uncertainties then cancel. Thus we write

$$
R = \frac{\sigma_W^{\ell\nu}}{\sigma_Z^{\ell\ell}}
$$

(2.5.2)

$$
= \frac{\sigma_W \Gamma_W^{\ell\nu} \Gamma_Z^{\text{tot}}}{\sigma_Z \Gamma_Z^{\ell\ell} \Gamma_W^{\text{tot}}} \equiv R_\sigma R_T \frac{\Gamma_Z^{\text{tot}}}{\Gamma_W^{\text{tot}}}
$$

where we have introduced a shorter notation for the partial leptonic and total cross sections. This relation has traditionally been used for ‘neutrino counting’ in hadron collider experiments\(^{39}\), which essentially amounted to an indirect determination of $\Gamma_Z^{\text{tot}}$ from the measured value of $R$ and the SM prediction for $\Gamma_W^{\text{tot}}$. More recently, it was pointed out\(^{40}\) that with $\Gamma_Z^{\text{tot}}$ measured by LEP the argument could be reversed yielding a determination of $\Gamma_W^{\text{tot}}$. In order to do so, both theoretical and experimental input is needed.

The ratio of the total cross sections, $R_\sigma = \sigma_W/\sigma_Z$, can be calculated in QCD. The uncertainties on $\sigma_W$ and $\sigma_Z$ associated with higher order corrections cancel almost completely in the ratio\(^{41}\). The main uncertainty on $R_\sigma$ comes from the choice of the structure functions. This choice is constrained, however, by recent measurements\(^{42,43}\) of the ratio of structure
functions \( F^W_2/F^Z_2 \), which reduces the uncertainty on \( R_{\sigma} \) to about 2%. The ratio of the partial widths, \( R_{\Gamma} \equiv \Gamma^W_{\ell \nu}/\Gamma^Z_{\ell \ell} \), can be expressed in terms of \( M_W/M_Z \) and \( \sin^2\theta_W \). The ratio of the vector boson masses has been measured in experiments at hadron colliders\(^{44,45} \) and can be inferred from results from deep inelastic scattering experiments\(^{46,47} \) while \( \sin^2\theta_W \) is determined by \( M_Z \). With this input, \( R_{\Gamma} \) is known with an uncertainty of about 1%.

The ratio \( R \) can be expressed in terms of experimentally measured quantities as follows:

\[
R = \frac{\sigma^W_{\ell \nu}}{\sigma^Z_{\ell \ell}} = \frac{N_W - N_W^{bg}}{\varepsilon_W L_W} \frac{\varepsilon_Z L_Z}{N_Z - N_Z^{bg}},
\]

(2.5.3)

where \( N_W \) and \( N_Z \) are the number of \( W \) and \( Z \) events observed, \( N_W^{bg} \) and \( N_Z^{bg} \) the number of background events, \( \varepsilon_W \) and \( \varepsilon_Z \) the overall detection efficiencies and \( L_W \) and \( L_Z \) the integrated luminosities. Usually, the two luminosities are equal and therefore cancel. Many systematic uncertainties that contribute to the individual efficiencies also cancel in the ratio. At present, the experimental error on \( R \) is dominated by the statistical error on the number of \( Z^0 \)'s.

In chapters three and four of this thesis, we will consider in more detail the theoretical input which enters relation (2.5.2), while chapters five, six and seven are dedicated to the experimental determination of the quantities on the rhs of equation (2.5.3) from recent UA1 data. In chapter eight finally, the measured value of \( R \) and the inferred value of \( \Gamma^W_{\ell \nu} \) are presented. Similar analyses, to be discussed in the next section, have been reported by the UA2 and CDF collaborations. Since the experimental error on \( R \) is dominated by the statistical error on the number of \( Z^0 \)'s and not by systematics, we combine all available measurements of \( R \) to reduce the error on \( \Gamma^W_{\ell \nu} \).

2.6 Determinations of \( \Gamma^W_{\ell \nu} \) by Other Experiments

2.6.1 Introduction

So far, three attempts have been made to determine \( \Gamma^W_{\ell \nu} \). The first concerns a direct measurement of the UA1 collaboration\(^8 \) based on old data taken in the 1983—1985 collider runs. The other two have only recently been reported by the UA2\(^22 \) and CDF\(^48 \) collaborations. They employ the same method as outlined in the previous section. We briefly discuss the three results, the UA1 result simply to contrast the difference in sensitivity obtained with the direct and indirect method, the other two because they will be reconsidered in chapter eight.

2.6.2 UA1, a Direct Measurement

The measurement consists of a (partial) reconstruction of the mass and width of the \( W \) resonance in \( W \rightarrow e\nu \) events. The data sample corresponds to an integrated luminosity of 0.77 pb\(^{-1} \) and was taken at \( \sqrt{s} = 0.546 \text{ TeV} \) and \( \sqrt{s} = 0.630 \text{ TeV} \). In UA1 (in the period up to 1985), electrons are identified by the presence of a high momentum charged track in the
Central Driftchamber matched in position and momentum by an isolated cluster in the electromagnetic (EM) calorimeter. The presence of energetic neutrinos is inferred from an apparent energy imbalance in the plane transverse to the beam (missing energy) in both the EM and hadronic calorimeters (see section 6.2.3). The EM calorimeter covers an angular range of: $5^\circ < \theta < 175^\circ$ with respect to the beam and has an energy resolution of: $15\%/E$ in the central region and $12\%/E$ in the forward region. A description of the other detector components can be found in chapter five.

Basically, the W selection requires isolated electrons, while both the electron and neutrino must have a transverse energy ($E_T$) larger than 30 GeV (this is much tighter than the standard selection which requires $E_T > 15$ GeV, but ensures a practically background-free sample). With these cuts they select 149 events. Unfortunately, the W mass cannot be completely reconstructed, because the longitudinal component of the neutrino momentum is not measured. Instead one uses a variable based on transverse momentum components only, called the transverse mass, $M_T$. It is defined as

$$M_T = \sqrt{2 p_T^e p_T^\nu (1 - \cos \phi_{e\nu})},$$

(2.6.1)

where $p_T^e$ and $p_T^\nu$ are the charged lepton $\{e, \mu, \tau\}$ and neutrino transverse momenta, and $\phi_{e\nu}$ is the relative azimuthal angle between the leptons. The shape of the distribution depends on the mass and width of the W boson, on the W momentum at production, and on the experimental selection biases and resolution. The $M_T$ distribution for the 149 events observed in the data is shown in figure 2.10.
2. Motivation for the Determination of $\Gamma_{W}^{\text{tot}}$

A Monte Carlo simulation of W production including all known acceptance and resolution effects was used to generate $M_{T}$ distributions. To determine $M_{W}$, sets of Monte Carlo events were generated with different $M_{W}$ values and with $\Gamma_{W}^{\text{tot}}$ fixed to 2.8 GeV. The result of a maximum likelihood fit on the 149 $W \rightarrow e\nu$ events gave:

$$M_{W} = 82.7 \pm 1.0 \text{ (stat)} \pm 2.7 \text{ (syst)} \text{ GeV.}$$

(2.6.2)

The $M_{T}$ distribution for the Monte Carlo events with the most likely value for $M_{W}$ has been indicated in figure 2.10 by the solid line. Additional sets of Monte Carlo events were generated for different values of $\Gamma_{W}^{\text{tot}}$ at the fitted W mass. A fit to the experimental $M_{T}$ distribution yielded as most probable value

$$\Gamma_{W}^{\text{tot}} = 2.8^{+1.4}_{-1.5} \text{ (stat)} \pm 1.3 \text{ (syst)} \text{ GeV.}$$

(2.6.3)

The 90% C.L. limit is $\Gamma_{W}^{\text{tot}} \leq 5.4$ GeV.

2.6.3 UA2, $\Gamma_{W}^{\text{tot}}$ from the Ratio $R$

The UA2 analysis is based on electron data corresponding to an integrated luminosity of 7.8 pb$^{-1}$ taken in the 1988 and 1989 runs at $\sqrt{s} = 0.63$ TeV. The UA2 detector$^{49}$ had been rebuilt in the preceding period to enhance its electron and neutrino identification capabilities. Basically, the apparatus consists of a lead-scintillator (EM) and iron-scintillator (hadronic) calorimeter which covers the full azimuthal angle and extends down to a 6$^{o}$ polar angle from the beam pipe (or: $|\eta| < 3$, where $\eta = -\ln(\tan(\theta/2))$ is the pseudorapidity). The room inside the calorimeter is occupied by a central detector which consists of three devices: an inner tracking detector (ITD), a transition radiation detector and a scintillating fibre detector (SFD). The latter contains a preshower detector to localize the early development of EM showers. Charged tracks and the position of the vertex are reconstructed using the SFD in conjunction with the ITD. Electrons are identified by an isolated cluster in the EM calorimeter matched in position by a charged track. Since UA2 has no magnetic field, momentum-energy matching is not possible. The separation of electrons from charged hadrons is achieved with the TRD. Neutrinos are reconstructed as in UA1.

$W \rightarrow e\nu$ events have been selected by the following kinematical cuts: (i) $p_{T}^{\ell} > 20$ GeV, (ii) $p_{T}^{\nu} > 20$ GeV and (iii) $M_{T} > 40$ GeV. In total 2041 events passed the selection, for which the QCD background (from overlapping $\pi^{0}$ and $\pi^{\pm}$ tracks) was found to be negligible. The only serious background comes from: $W \rightarrow \tau\nu_{\tau}$ followed by $\tau \rightarrow e\nu_{e}\nu_{\tau}$. Its contribution was calculated by Monte Carlo, which yielded 75.7 events.

$Z \rightarrow e^{+}e^{-}(\gamma)$ candidates have been selected from an event sample containing at least two EM clusters with an invariant mass larger than 40 GeV (if a third cluster with transverse energy exceeding 5 GeV was present, it was included in the invariant mass calculation). The observed invariant mass spectrum consists of an exponentially falling distribution associated with the Drell-Yan continuum and a QCD background from two-jet events, with a Breit-Wigner from the $Z^{0}$ resonance superimposed. The final $Z$ selection required an invariant mass larger than 76 GeV. In total 169 events fulfilled this criterion. The QCD background (2.4 events) was obtained by extrapolation of the exponential from the low mass
region, while the fraction of Drell-Yan events (1.65% from single photon exchange and $\gamma^*Z$ interference) was computed by Monte Carlo.

The acceptance calculation was spread over three calorimetric regions yielding varying results over the different regions. Applying relation (2.5.3) they find for the ratio $R$

$$R = 9.38^{+0.82}_{-0.72} \text{ (stat)} \pm 0.25 \text{ (syst).} \quad (2.6.4)$$

In order to determine $\Gamma_W^{\text{tot}}$ via relation (2.5.2), they use their own measurement of the ratio of the vector boson masses: $M_W/M_Z = 0.8831 \pm 0.0048 \pm 0.0026$ and the LEP and SLC measurement of the Z mass: $M_Z = 91.15$ GeV to fix the W mass to $M_W = 80.5$ GeV and the weak mixing parameter to $\sin^2\theta_W = 0.220$. This yields for the ratio of the partial leptonic widths: $R_\tau = 2.717$. For the ratio of the total cross sections the value $R_\sigma = 3.116$ was used, while for the total width of the $Z^0$ a weighted average of LEP and SLC results was taken: $\Gamma_Z^{\text{tot}} = 2.546 \pm 0.032$ GeV. They then obtain the following value for the total W width

$$\Gamma_W^{\text{tot}} = 2.30 \pm 0.19 \text{ (stat)} \pm 0.06 \text{ (syst)} \text{ GeV.} \quad (2.6.5)$$

### 2.6.4 CDF, $\Gamma_W^{\text{tot}}$ from the Ratio $R$

Also the CDF measurement is based on electron events. The data sample was taken at $\sqrt{s} = 1.8$ TeV and corresponds to an integrated luminosity of 4.4 pb$^{-1}$. The CDF detector consists of central tracking chambers enclosed in a solenoidal magnet (with a maximum field of 1.5 T allowing to measure track momenta with a resolution of: $\Delta p_T^2/p_T^2 = 0.001$ GeV$^{-1}$) surrounded by calorimetry and muon spectrometers. The EM and hadronic calorimeters, which use lead and iron as absorber respectively, cover the full azimuthal angle and extend down to a 1.7° polar angle from the beam pipe ($|\eta| < 4.2$). They are split into three different regions: the central ($|\eta| < 1.1$), plug ($1.1 < |\eta| < 2.4$) and forward ($2.4 < |\eta| < 4.2$) regions. Only the central calorimeters employ scintillator as a sampling medium, the others use gas-proportional tubes. Electron and neutrino identification proceeds basically in the same manner as described above for UA1.

W and Z candidates were selected from a common sample of events with at least one well measured, isolated, high $p_T$ electron in the central EM calorimeter, satisfying: $E_T > 20$ GeV and $|\eta| < 1.0$. It was also required that there be no additional clusters with transverse energy larger than 10 GeV other than the electron(s) in the event (this ‘zero jet’ requirement reduces systematic uncertainties and backgrounds).

The W selection further required a neutrino with $E_T > 20$ GeV. In total 1828 events passed the selection. The background, computed by Monte Carlo, consists of 101 events and is dominated by $W \rightarrow \tau\nu$ decays.

Z candidates were selected by requiring a second EM cluster with $E_T > 20$ GeV not necessarily in the central calorimeter, but away from the edges. The invariant mass of the two clusters should lie within 65 GeV and 115 GeV. These criteria were satisfied by 193 events for which the background was computed to be 6 events (mainly QCD background).
2. Motivation for the Determination of $\Gamma_w^{\text{tot}}$

Again, the acceptance calculation is divided into three calorimetric regions. From a global ratio of $Z \rightarrow e^+e^-$ and $W \rightarrow e\nu$ detection efficiencies of: $\varepsilon_Z/\varepsilon_W = 1.04 \pm 0.03$, and the observed numbers of events the following value for the ratio $R$ has been derived

$$R = 10.2 \pm 0.8 \, \text{(stat)} \pm 0.4 \, \text{(syst)}.$$  \hspace{1cm} (2.6.6)

This value has been corrected for the 'zero jet' requirement which is expected to increase the ratio $R$ by $0.8 \pm 0.5\%$. In order to determine $\Gamma_w^{\text{tot}}$, the following input values have been used$^{48}$: $\sin^2\theta_W = 0.229$, $R_\sigma = 3.23 \pm 0.03$, $R_\Gamma = 2.70 \pm 0.02$ and $\Gamma_Z^{\text{tot}} = 2.57 \pm 0.07$ GeV. The measured value of $R$ then leads to the following value for the $W$ total width

$$\Gamma_w^{\text{tot}} = 2.19 \pm 0.20 \, \text{GeV}.$$  \hspace{1cm} (2.6.7)
3 \ W^\pm \ and \ Z^0 \ Decays
3. $W^\pm$ and $Z^0$ Decays
3.1 Introduction

In this chapter, SM predictions for the partial and total widths of the weak IVBs are considered. Our main interest concerns the experimental and theoretical uncertainties on \( R_\Gamma \) and \( \Gamma_\text{tot}^{\text{W}} \). In lowest order (the Born approximation), the partial widths can be expressed in terms of only a few electroweak parameters, namely: the IVB masses, \( G_F \) and \( \sin^2 \theta_W \). All of these have been well measured, the largest uncertainties being on \( M_W \) and \( \sin^2 \theta_W \). Since \( M_W \) enters the W width as a third power, while \( \sin^2 \theta_W \) also has an impact on \( R_\sigma \) (to be discussed in the next chapter), it is worthwhile to optimize the available experimental information on these quantities. Higher order corrections introduce a dependence on other parameters of the theory as well. An improved Born approximation exists, which on the one hand contains all ‘large’ corrections, but on the other hand retains the simple form of the lowest order expression\(^{51}\). For decays into heavy fermions, an additional phase space factor is required, while QCD corrections become mass dependent as well.

In the following, we first give the Born level expressions for the partial widths and then deduce a value for \( M_W \) from measurements of the IVB masses and \( R_\sigma \) (see also section 2.3.3). Along the way, we get a rather accurate determination of \( \sin^2 \theta_W \). However, we will consider another definition of this quantity in the subsequent section which allows to decrease the uncertainty still further by a factor two. We then give the improved Born approximation with numerical predictions for the decay widths. After discussing the effect on the widths from decays into heavy quarks, we dedicate one section to IVB decays into charginos and neutralinos.

3.2 \( W^\pm \) and \( Z^0 \) Partial Widths in the Born Approximation

In the Born approximation, the partial widths associated with W and Z decays into massless fermion pairs: \( W \rightarrow f\bar{f}' \) and \( Z^0 \rightarrow f\bar{f} \), are given by

\[
\Gamma_W^{ff'} = N\frac{G_FM_W^3}{6\pi\sqrt{2}}
\]

\[
\Gamma_Z^{ff} = N\rho_0\frac{G_FM_Z^3}{24\pi\sqrt{2}} \left[ 1 + (1 - 4|Q_f|\sin^2 \theta_W)^2 \right],
\]

where \( N = 1 \) for leptons, \( N = 3 \) for quarks and \( Q_f \) is the fermion charge. In the minimal SM the rho-parameter \( \rho_0 \) is equal to 1. In the case of W decays into quark-antiquark pairs, formula (3.2.1) refers to the sum over all down-like quarks associated with a given up-like quark. For a given down-like quark, \( q' \), a factor \( |V_{qf}|^2 \) would appear, where \( V_{qf} \) is the relevant term of the Cabibbo-Kobayashi-Maskawa matrix (\( \Sigma_{qf}^1|V_{qf}|^2 = 1 \) by unitarity).

It is interesting to note that the quantity in square brackets in relation (3.2.2) which leads to a \( \sin^2 \theta_W \)-dependence of \( R_\Gamma \) exhibits a minimum in \( \sin^2 \theta_W = 0.25 \) for charged leptons (\( Q_\ell = 1 \)). Since \( \sin^2 \theta_W \) has been experimentally determined to be approximately 0.23, this means that \( R_\Gamma \) (which is a ratio of leptonic widths) is relatively insensitive to variations in \( \sin^2 \theta_W \). The dependence is stronger for the coupling of a \( Z^0 \) to quarks (therefore, variations in \( \sin^2 \theta_W \) will have a bigger effect on \( R_\sigma \)).
Table 3.1: Measurements of the weak IVB masses. The first error is statistical, the second systematic and the numbers in italics are uncertainties on the absolute energy scale. In taking the average, all errors have been added in quadrature.

<table>
<thead>
<tr>
<th>Experiment</th>
<th>MW (GeV)</th>
<th>MZ (GeV)</th>
<th>MW/MZ</th>
</tr>
</thead>
<tbody>
<tr>
<td>UA1 (e)8</td>
<td>82.7 ± 1.0 ± 2.7</td>
<td>93.1 ± 1.0 ± 3.1</td>
<td>0.888 ± 0.014</td>
</tr>
<tr>
<td>UA2 (e)44</td>
<td>80.79 ± 0.31 ± 0.21 ± 0.81</td>
<td>91.49 ± 0.35 ± 0.12 ± 0.91</td>
<td>0.8830 ± 0.0048 ± 0.0026</td>
</tr>
<tr>
<td>CDF (e)52</td>
<td>79.91 ± 0.35 ± 0.24 ± 0.19</td>
<td>91.37 ± 0.34 ± 0.04 ± 0.22</td>
<td>0.8746 ± 0.0050 ± 0.0027</td>
</tr>
<tr>
<td>CDF (μ)52</td>
<td>79.90 ± 0.53 ± 0.32 ± 0.08</td>
<td>90.71 ± 0.40 ± 0.04 ± 0.09</td>
<td>0.8808 ± 0.0070 ± 0.0035</td>
</tr>
<tr>
<td>average</td>
<td>80.08 ± 0.34</td>
<td>91.09 ± 0.28</td>
<td>0.8799 ± 0.0034</td>
</tr>
<tr>
<td>MARK II9</td>
<td></td>
<td>91.14 ± 0.12</td>
<td></td>
</tr>
<tr>
<td>LEP9</td>
<td></td>
<td>91.172 ± 0.031</td>
<td></td>
</tr>
</tbody>
</table>

Measurements of GF, MW and MZ are sufficient to compute the expressions (3.2.1) and (3.2.2), because at tree level sin²θW is related to the IVB masses by

\[
\sin^2 \theta_W = 1 - \frac{M_W^2}{M_Z^2}.
\]  

(3.2.3)

Table 3.1 shows measurements of the weak IVB masses coming from experiments at pp and e⁺e⁻ colliders. The outstanding result, of course, is the measurement of MZ by the LEP experiments with an uncertainty of only 0.3%. However, recent measurements by the UA2 and CDF collaborations have also substantially reduced the uncertainty on MW. Averaging the UA1, UA2 and CDF results while combining statistical and systematic errors in quadrature yields a relative uncertainty on MW of 0.4%. This uncertainty can be further reduced by virtue of relation (3.2.3): since MZ has been measured with high precision, a measurement of sin²θW yields a determination of MW and vice versa. This will be considered in the next section.

3.3 Determination of MW from Measurements of MW, MZ and RV

Direct measurements of MW (and MZ) at hadron colliders suffer from relatively large uncertainties on the absolute energy scale. As can be seen from table 3.1, they are even dominant for UA1 and UA2. In the measurement of the ratio MW/MZ (last column of table 3.1), this uncertainty drops out. Therefore, a better determination of MW might be obtained by combining the measurement of the mass ratio with the measurement of MZ from LEP, yielding a rescaled value for MW. Direct measurements of MW are compared to the rescaled values in table 3.2. Statistical and systematic errors have been added in quadrature.
Table 3.2: $M_W$ (GeV) measured by $p\bar{p}$ collider experiments and rescaled to $M_Z$ from LEP. Errors have been added in quadrature.

<table>
<thead>
<tr>
<th></th>
<th>UA1 (e)</th>
<th>UA2 (e)</th>
<th>CDF (e)</th>
<th>CDF ((\mu))</th>
</tr>
</thead>
<tbody>
<tr>
<td>$M_W$ measured</td>
<td>82.7 ± 2.9</td>
<td>80.79 ± 0.89</td>
<td>79.91 ± 0.47</td>
<td>79.90 ± 0.62</td>
</tr>
<tr>
<td>$M_W$ rescaled</td>
<td>81.0 ± 1.3</td>
<td>80.50 ± 0.50</td>
<td>79.74 ± 0.52</td>
<td>80.30 ± 0.71</td>
</tr>
</tbody>
</table>

For UA1 and UA2 the gain is clear, while for CDF the result degrades. This can be easily understood, because the scale uncertainty for CDF is not that large and renormalization of $M_Z$ to the LEP result transfers the statistical error associated with a relatively small $Z^0$ sample onto $M_W$. Therefore, the best determination of $M_W$ is obtained by taking a weighted average of the rescaled values of the UA1 and UA2 results and the directly measured values of CDF, yielding

$$
(M_W)_{M_W,M_Z} = 80.16 ± 0.29 \text{ GeV.}
$$

(3.3.1)

Combining this value with $M_Z$ from LEP gives the following value for the weak mixing parameter

$$
(sin^2\theta_W)_{M_W,M_Z} = 0.2270 ± 0.0056,
$$

(3.3.2)

which is free of theoretical uncertainties. The error is dominated by low statistics of the $Z^0$ samples at hadron colliders.

Another way to obtain $M_W$ is to combine values of $sin^2\theta_W$ extracted from neutral current data (section 2.3.3) with $M_Z$. In many cases, however, the value extracted depends strongly on $m_t$ through radiative corrections. An exception is formed by the determination of $sin^2\theta_W$ from $R_V$, which is almost independent of $m_t$ (see figure 2.7 in chapter 2). The most precise measurements come from the CHARM$^{31}$ and CDHS$^{32}$ collaborations at CERN

CHARM: \[ (sin^2\theta_W)_{R_V} = 0.236 ± 0.007, \]

CDHS: \[ (sin^2\theta_W)_{R_V} = 0.228 ± 0.007. \]

(3.3.3)

Combining these measurements with the value in (3.3.2) gives

$$
sin^2\theta_W = 0.230 ± 0.004,
$$

(3.3.4)

which together with $M_Z$ from LEP yields for the W mass

$$
M_W = 80.00 ± 0.21 \text{ GeV.}
$$

(3.3.5)

These two values are practically independent of the t-quark mass and do not require higher order corrections.
3. W$^\pm$ and Z$^0$ Decays

3.4 Determination of $\sin^2 \theta_W$ from $M_Z$

We take advantage of the precise $M_Z$ measurement from LEP by extracting $\sin^2 \theta_W$ from the following tree-level relation

$$\sin^2 \theta_W \cos^2 \theta_W = \frac{\pi \alpha}{\sqrt{2} G_F M_Z^2},$$

(3.4.1)

for which the experimental uncertainties on the quantities on the rhs are extremely small. The main source of uncertainty originates from radiative corrections $\Delta r$ which modify this relation by

$$\sin^2 \theta_W \cos^2 \theta_W = \frac{\pi \alpha}{\sqrt{2} G_F M_Z^2 (1 - \Delta r)}.$$

(3.4.2)

In this section, we show that the theoretical uncertainty introduced by $\Delta r$ is small enough to make this method of determination of $\sin^2 \theta_W$ more attractive than the one employing the IVB masses and $R_V$.

The quantity $\Delta r$ depends on all parameters of the SM, in particular on the unknown $m_t$ and $m_h$. Its precise form depends on the details of a chosen renormalization scheme. One scheme, already mentioned in section 2.3.3, defines relation (3.2.3) to remain valid in all orders. This leads to a strong $m_t$-dependence of $\sin^2 \theta_W$ for $m_t$ larger than about 60 GeV (see the curve labelled VBM in figure 2.7). Here we seek to minimize such a dependence. This can be achieved by defining $\sin^2 \theta_W = \tilde{s}_W^2$ as the effective weak mixing angle that occurs in the amplitude for the coupling of fermions to an on-shell $Z^0$ boson. In that case, the radiative corrections can be approximated by

$$\Delta r \approx \delta \alpha - \delta \rho,$$

(3.4.3)

where $\delta \alpha$ accounts for the running of $\alpha$ from the electron mass up to the $Z^0$ mass

$$\frac{\alpha(M_Z)}{\alpha} = \frac{1}{1 - \delta \alpha},$$

(3.4.4)

and is given by

$$\delta \alpha = 0.0601 + \frac{40}{9} \frac{\alpha}{\pi} \ln \left( \frac{M_Z}{91 \text{ GeV}} \right) \pm 0.0009,$$

(3.4.5)

and $\delta \rho$ consists of two pieces

$$\delta \rho = \delta \rho_t + \delta \rho_h,$$

(3.4.6)

one depending quadratically on $m_t$

$$\delta \rho_t = \frac{3 G_F m_t^2}{8 \pi^2 \sqrt{2}} + \frac{\left( \frac{3 G_F m_t^2}{8 \pi^2 \sqrt{2}} \right)^2}{3} \frac{19 - 2 \pi^2}{3} + O[(G_F m_t^2)^3],$$

(3.4.7)

and the other depending logarithmically on $m_h$

$$\delta \rho_h = -\frac{11 G_F m_W^2}{24 \pi^2 \sqrt{2}} \tan^2 \theta_W \ln \left( \frac{m_h^2}{M_W^2} \right).$$

(3.4.8)
3.4 Determination of $\sin^2\theta_W$ from $M_Z$

![Figure 3.1: The $m_t$-dependence of $\sin^2\theta_W$.](image)

By using equations (3.4.3) and (3.4.4), while absorbing $\delta\alpha$ in $\alpha(M_Z)$, one can rewrite equation (3.4.2) as follows

$$\sin^2\theta_W \cos^2\theta_W = \frac{\pi \alpha(M_Z)}{\sqrt{2} G_F} \frac{1}{\rho M_Z^2}, \quad (3.4.9)$$

where $\rho \approx 1 + \delta\rho$. For $m_t \in [50, 200]$ GeV and $m_h \in [50, 1000]$ GeV the value of $\rho$ is restricted to $\rho = 1.005 \pm 0.006$ and the value of $\sin^2\theta_W$ to

$$\sin^2\theta_W = 0.230 \pm 0.002. \quad (3.4.10)$$

The individual uncertainties ($\delta\sin^2\theta_W = 0.0002$ (M_Z), 0.0007 (m_h) and 0.0020 (m_t)) have been added in quadrature. Figure 3.1 shows the $m_t$-dependence of $\sin^2\theta_W$ where the shaded band reflects the uncertainty on $m_h$. If we assume the t-quark to be light ($\sim 50$ GeV), then the weak mixing parameter becomes

$$\sin^2\theta_W = 0.231 \pm 0.001, \quad \text{for } m_t = 50 \text{ GeV}, \quad (3.4.11)$$

where the only uncertainty now comes from $M_Z$ and $m_h$. In that case, the rho-parameter is: $\rho = 0.999$.  

3. W\pm and Z^0 Decays

3.5 Numerical Predictions for W\pm and Z^0 Decay Widths in the SM

3.5.1 Improved Born Approximation

Higher order corrections not only modify relations between parameters (like e.g. relation (3.4.1)), they also modify the predictions for the observables themselves. The Born approximation can be rewritten such, that it comprises all 'large' corrections

\[ \Gamma_W^{\text{ff}} = N \frac{G_F M_W^3}{6\pi \sqrt{2}} \left[ 1 + \delta_W^f \right] \]  
(3.5.1)

\[ \Gamma_Z^{\text{ff}} = N \rho \frac{G_F M_Z^3}{24\pi \sqrt{2}} \left[ 1 + \left( 1 - 4|Q_f|\bar{s}_W^2 \right)^2 \right] \left[ 1 + \delta_Z^f \right], \]  
(3.5.2)

with \( \rho \) and \( \bar{s}_W^2 \) as defined in the previous section, and where

\[ N = 1 \quad \text{for leptons,} \]  
(3.5.3)

\[ N = 3 \left( 1 + \frac{\alpha_S}{\pi} \right) \quad \text{for quarks.} \]

The last factor in (3.5.3) comes from QCD corrections. A recent global analysis of experimental results on the strong coupling constant yields for \( \alpha_S \) at the \( Z^0 \) mass: \( \alpha_S(M_Z) = 0.11 \pm 0.01 \)\(^{56} \). From this we obtain a value for \( \alpha_S/\pi \) restricted to

\[ \frac{\alpha_S}{\pi} = 0.035 \pm 0.003 \]  
(3.5.4)

The electroweak radiative corrections, \( \delta_W^f \) and \( \delta_Z^f \), turn out to be rather small (\( \delta^f < 0.1\% \)) and for the purpose of this thesis can be neglected\(^{57} \).

In table 3.3, numerical predictions for the \( W\pm \) and \( Z^0 \) decay widths are given, computed in the improved Born approximation (phase space factors have been included, yielding slightly reduced widths for c- and b-quarks). The individual uncertainties on the total widths are for the \( W\pm \): \( \delta\Gamma_{W\pm}^{\text{tot}} = 17 \text{ MeV} (M_W) \) and 4 MeV (\( \alpha_S \)), and for the \( Z^0 \): \( \delta\Gamma_{Z}^{\text{tot}} = 3 \text{ MeV} (M_Z) \), 5 MeV (\( \alpha_S \)) and 20 MeV (\( m_t \) and \( m_b \)). They have been added in quadrature.

Now we can also compute \( R_{\Gamma} \) which is given by

\[ R_{\Gamma} = 2.672 \pm 0.030. \]  
(3.5.5)

The uncertainty is dominated by the \( W \) mass (\( \delta R_{\Gamma} = 0.021 \)) and by the unknown t-quark mass via \( \bar{s}_W^2 \) and \( \rho \) (\( \delta R_{\Gamma} = 0.021 \)). If we assume the t-quark to be light, the latter uncertainty disappears and we get

\[ R_{\Gamma} = 2.690 \pm 0.021, \quad \text{for } m_t = 50 \text{ GeV}. \]  
(3.5.6)
3.5 Numerical Predictions for $W^\pm$ and $Z^0$ Decay Widths in the SM

Table 3.3: SM predictions for $W^\pm$ and $Z^0$ decay widths in the improved Born approximation with: $M_W = 80.00 \pm 0.21$ GeV, $M_Z = 91.172 \pm 0.031$ GeV, $\alpha_S = 0.11 \pm 0.01$, $\delta_W = 0.230 \pm 0.002$ and $\rho = 1.005 \pm 0.006$. Uncertainties have been indicated in brackets.

<table>
<thead>
<tr>
<th>W Decay Mode</th>
<th>Partial Width (MeV)</th>
<th>Z Decay Mode</th>
<th>Partial Width (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$e\nu_e$, $\mu\nu_\mu$, $\tau\nu_\tau$</td>
<td>224 (2)</td>
<td>$\nu_e\bar{\nu}<em>e$, $\nu</em>\mu\bar{\nu}<em>\mu$, $\nu</em>\tau\bar{\nu}_\tau$</td>
<td>167 (1)</td>
</tr>
<tr>
<td>$u(d+s+b)$</td>
<td>696 (6)</td>
<td>$e^+e^-$, $\mu^+\mu^-$, $\tau^+\tau^-$</td>
<td>84 (1)</td>
</tr>
<tr>
<td>$c(d+s+b)$</td>
<td>695 (6)</td>
<td>$u\bar{u}$, $c\bar{c}$</td>
<td>297 (3)</td>
</tr>
<tr>
<td>Total</td>
<td>2063 (17)</td>
<td>$d\bar{d}$, $s\bar{s}$</td>
<td>383 (3)</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$b\bar{b}$</td>
<td>379 (3)</td>
</tr>
<tr>
<td>Total</td>
<td>2492 (20)</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

3.5.2 Decays into Massive Fermions

Expressions (3.5.1) and (3.5.2) for the partial widths assume the daughter particles of $W$ and $Z$ decays to be massless, which is a very good approximation for the decays considered so far. However, for non-negligible masses the expressions need to be modified to account for phase space suppression. If the masses of the two fermions of a $W$ decay are given by $m_a$ and $m_b$, and for a $Z$ decay by $m$, then the widths are given by\(^58\)

$$\Gamma_{W}^{ff} = \frac{N G_F M_W^3}{6\pi \sqrt{2}} \left[ 1 - \frac{1}{2} \frac{m_a^2}{M_W^2} + \frac{1}{2} \frac{m_b^2}{M_W^2} \right],$$

$$\Gamma_{Z}^{ff} = \frac{N_{\rho} G_F M_Z^3}{24\pi \sqrt{2}} \left[ 1 - 4x + (1 - 4Q_{f}\delta_W^2)^2(1 + 2x) \right], \quad (3.5.7)$$

where $x_a = m_a^2/M_W^2$, $x_b = m_b^2/M_W^2$, $x = m^2/M_Z^2$, and $\lambda(1, x_a, x_b)$ is defined in (2.4.3). These additional factors account for the slight reduction of partial widths involving c- and b-quarks in table 3.3.

For the decays: $W \to t\bar{b}$ and $Z^0 \to t\bar{t}$, one also has to take into account the $m_t$-dependence of the QCD corrections\(^59\). For $W$ decays, they increase from 3.5% to about 8% (with respect to the zero mass width) if $m_t$ goes from 0 to 40 GeV. For $Z$'s, the correction becomes almost 9% for top masses around 30 GeV. For larger $m_t$ it slowly vanishes, while close to the $W$ or $Z$ mass it starts to oscillate rapidly\(^60\).

Figure 3.2 gives an overview of the partial and total widths of $W$'s and $Z$'s in the SM. It illustrates the relative contributions of leptonic and hadronic partial widths and the size of QCD corrections. Also shown is the effect of phase space suppression on the width from a non-negligible $t$-quark mass, and the variation of the QCD correction with $m_t$. 
3.6 $W^\pm$ and $Z^0$ Decays into Charginos and Neutralinos in the MSSM

In chapter two, we mentioned that partial contributions to the $W^\pm$ and $Z^0$ widths from supersymmetric decays into charginos and neutralinos can become large enough to be detected with the present experimental resolution. An example was given for a specific choice of mass matrix parameters defining mixing in the chargino-neutralino sector, which allowed to write down a simple expression for the $W^\pm$ width in terms of the couplings and a phase space factor. Since this choice is not preferred over any other one, it is better to use more general expressions, which for completeness will be given here.
Mixing of winos and higgsinos into two chargino mass eigenstates, $\tilde{\chi}^\pm$ and $\tilde{\chi}'^\pm$, is described by the $2 \times 2$ matrix

$$
\begin{pmatrix}
M_2 & \sqrt{2}M_W \sin \beta \\
\sqrt{2}M_W \cos \beta & \mu
\end{pmatrix}
$$

(3.6.1)

where $M_2$ is the SU(2) gaugino mass, $\mu$ is a Higgs mixing term, and $\tan \beta \equiv \frac{v_2}{v_1}$ with $v_{1,2}$ the vacuum expectation values of the two Higgses $\tilde{H}^0_{1,2}$. The same parameters characterize the $4 \times 4$ neutralino mass matrix (in grand unified theories $M_1$ is related to $M_2$), which in the basis $\{B, \tilde{W}_3, \tilde{H}^0_1, \tilde{H}^0_2\}$ is given by

$$
\begin{pmatrix}
M_1 & 0 & -M_Z \cos \beta \sin \theta_W & M_Z \sin \beta \sin \theta_W \\
0 & M_2 & M_Z \cos \beta \cos \theta_W & -M_Z \sin \beta \cos \theta_W \\
-M_Z \cos \beta \sin \theta_W & M_Z \cos \beta \cos \theta_W & 0 & -\mu \\
M_Z \sin \beta \sin \theta_W & -M_Z \sin \beta \cos \theta_W & -\mu & 0
\end{pmatrix}
$$

(3.6.2)

It has been argued$^{36}$, that the three parameters should lie in the following ranges: $0 \leq M_2 \leq 200$ GeV, $0 \leq |\mu| \leq 250$ GeV and $1 \leq \tan \beta < 10$. The chargino and neutralino masses are obtained by diagonalization of these matrices.

The partial width for the decay, $W^\pm \to \tilde{\chi}^\pm \tilde{\chi}_i$, was computed in terms of the parameters $(M_2, \mu, \tan \beta)$ in 1983 assuming massless neutralinos$^{61}$. Recently, the formula has been generalized to account for non-zero neutralino masses. Normalized to the partial width for $W^\pm \to ev$, it is given by$^{36}

$$
\Gamma_{W}^{\chi^\pm \chi_i} = 3 \Gamma_{W}^{ev} \lambda^{1/2}(1,x,y_i) \times

\left\{ \left[ 2 - x - y_i - (x - y_i)^2 \right] \frac{O_L^2 + O_R^2}{3} - 4O_L O_R (xy_i)^{1/2} \right\},
$$

(3.6.3)

where $x = m_{\tilde{\chi}}^2 / M_W^2$, $y_i = m_{\tilde{\chi}_i}^2 / M_W^2$ and $\lambda(1,x,y_i)$ has been defined in (2.4.3). The matrices $O_L$ and $O_R$ are defined in terms of other matrices which diagonalize (3.6.1) and (3.6.2)$^{18}$. The decay width for $Z^0 \to \tilde{\chi}^+ \tilde{\chi}^-$ is given by$^{61}

$$
\Gamma_{Z}^{\tilde{\chi}^+ \tilde{\chi}^-} = 4 \Gamma_{Z}^{\nu \nu} \left(1 - 4z\right)^{1/2} \left\{ \left[ 1 - z \right] \left( O_L^2 + O_R^2 \right) - 6O_L O_R \right\},
$$

(3.6.4)

where the primed matrices are defined in terms of ones diagonalizing (3.6.1). The partial width associated with a $Z^0$ decay into a pair of equal mass neutralinos can be obtained from (3.6.4) by insertion of a factor 1/2, because the final state particles are identical, and by replacing the primed matrices by double primed ones which diagonalize (3.6.2).

Figure 3.3 shows the sum of all contributions to the $W^\pm$ and $Z^0$ decay widths in $(M_2, \mu)$-parameter space for $\tan \beta = 2$. The maximum, attained for small $M_2$ and large $\mu$, is approximately 800 MeV for $W$'s and 700 MeV for $Z$'s. In this limit, the lightest neutralino is the photino$^{36}$. Since both surfaces look rather similar, it will be difficult to find regions in parameter space with substantially larger contributions to $\Gamma_W^{\text{tot}}$ than to $\Gamma_Z^{\text{tot}}$. The same figure shows the correlation between the widths and the masses of the lightest chargino and neutralino. For mass combinations above the kinematic limit, the widths become zero.
3. $W^\pm$ and $Z^0$ Decays

Figure 3.3: Contributions to the $W^\pm$ (a) and $Z^0$ (b) widths from decays into charginos and neutralinos, and the mass of the lightest chargino (c) and neutralino (d) in the $(M_2, \mu)$-plane for $\tan \beta = 2$. 62
4 $W^\pm$ and $Z^0$ Production in $p\bar{p}$ Collisions
4. $W^\pm$ and $Z^0$ Production in $p\bar{p}$ Collisions
4.1 Introduction

In this chapter, we consider the SM predictions for $W^\pm$ and $Z^0$ production cross sections, $\sigma_W$ and $\sigma_Z$ respectively, and their ratio $R_\sigma$. The theoretical framework is based on the QCD improved parton model, which is generally applicable to high energy processes involving a hard interaction. It deals with the composite nature of protons and antiprotons and allows to compute observable cross sections in a systematic way. The formalism for IVB production in particular has been developed by Drell and Yan\textsuperscript{63} in 1970, and since then has been extended by other people\textsuperscript{64}.

4.2 The Parton Model

High energy scattering of protons and antiprotons can be conveniently described in the parton model in terms of hard interactions between proton and antiproton constituents or partons (quarks and gluons). In the naive model as originally envisaged by Feynman\textsuperscript{65}, the scattering event is defined with respect to an infinite-momentum frame in which the proton and antiproton are moving with high speed (the LAB frame at pp colliders). In such a frame, relativistic time dilation slows down the rate at which partons interact with one another, so that the hard interaction leading to the scattering event occurs on a time scale short compared to the scale which controls the evolution of the parton system. During the hard interaction the partons can then be treated as though they were effectively free. The model allows to express observable cross sections in terms of parton densities and subprocess cross sections associated with interactions at parton level. It neglects, however, the dynamical role of gluons as carriers of the strong force. This is accounted for in the QCD improved parton model.

Figure 4.1 illustrates the parton model picture of the inclusive production of a vector boson $V$ in $pp$ collisions

$$p + \bar{p} \to V + X.$$  \hfill (4.2.1)

A proton with four-momentum $P_1$ and antiproton with four-momentum $P_2$ yield a parton of type $i$ and $j$ respectively. The partons carry a fraction of their parent (anti)proton momentum $p_1 = x_1 P_1$ and $p_2 = x_2 P_2$ with $0 \leq x_1, x_2 \leq 1$, and upon scattering create an IVB $V$ and additional partons in the subprocess

$$i + j \to V + \text{`partons'},$$  \hfill (4.2.2)

for which the cross section is given by $\hat{\sigma}_{ij}$. The experimentally observable cross section, $\sigma$, for the complete process in (4.2.1) is obtained by multiplying the subprocess cross section $\hat{\sigma}_{ij}$ by the probability to have a parton of type $i$ in the proton with momentum fraction between $x_1$ and $x_1 + dx_1$: $f_i(x_1)dx_1$, and the probability to have a parton of type $j$ in the antiproton with momentum fraction between $x_2$ and $x_2 + dx_2$: $f_j(x_2)dx_2$, integrating over all momentum fractions and summing over all parton types

$$\sigma = \sum_{i,j} \left[ \int_0^1 dx_1 \int_0^1 dx_2 f_i(x_1)f_j(x_2)\hat{\sigma}_{ij} \right].$$  \hfill (4.2.3)
4. $W^\pm$ and $Z^0$ Production in $p\bar{p}$ Collisions

![Diagram](image)

Figure 4.1: The inclusive production of a vector boson $V$ in proton-antiproton collisions: $p + \bar{p} \rightarrow V + X$, via a hard parton subprocess: $i + j \rightarrow V + \text{ 'partons'}$.

Other partons in the proton and antiproton, which do not participate in the hard interaction, are merely 'spectators'. In the experiment, they are responsible for some additional activity in the detector known as the 'underlying event'. However, most particles originating from these spectators disappear forwardly and remain inside the beam pipe.

Before discussing the IVB total cross sections, we introduce a few kinematical variables. The external proton-antiproton and the internal parton-parton CM energy squared will be indicated respectively by

\[ s \equiv (P_1 + P_2)^2, \]
\[ \hat{s} \equiv (p_1 + p_2)^2. \]  

(4.2.4)

If the (anti)proton mass is neglected, they are related according to

\[ \hat{s} = x_1 x_2 s. \]  

(4.2.5)

Furthermore, we define a variable $\tau$ as

\[ \tau \equiv M^2/s, \]

(4.2.6)

where $M$ is the vector boson mass.

4.3 The Total $W^\pm$ and $Z^0$ Production Cross Sections

4.3.1 The Born Approximation

In the naive parton model the only subprocess contributing to IVB production is quark-antiquark annihilation: $q + \bar{q} \rightarrow V$ (see the Born graph on the right), because gluons do not couple to weak IVBs. The subprocess cross sections can be simply obtained using standard Feynman rules for the couplings of the bosons to quarks. Averaging over initial

![Born graph](image)

Figure 4.2: Born graph.
colour states and quark spins and summing over the three IVB helicity states, we have for instance for $W^+$ production by a $u$- and $\bar{d}$-quark

$$\hat{\sigma}(u + \bar{d} \rightarrow W^+) = \frac{2}{3} \pi g^2_W \cos^2 \theta_C \delta(\hat{s} - M^2_W),$$  

(4.3.1)

where $g_W$ is the dimensionless weak coupling ($g_W = \pi \alpha/2 \sin^2 \theta_W$) and $\theta_C$ is the Cabibbo angle (for the $W^-$ cross section all particles should be replaced by their respective antiparticles), and for $Z^0$ production by a $u$- and $\bar{u}$-quark

$$\hat{\sigma}(u + \bar{u} \rightarrow Z^0) = \frac{1}{6} \pi \frac{g^2_W}{\cos^2 \theta_W} n_u \delta(\hat{s} - M^2_Z),$$  

(4.3.2)

where $n_q$ depends on the charge of a quark $q$, and is given by

$$n_q = 1 + (1 - 4|Q_q| \sin^2 \theta_W)^2.$$  

(4.3.3)

The total cross section for the production of weak IVBs in $p\bar{p}$ collisions, $\sigma_V$, is obtained by applying equation (4.2.3) allowing for all possible $q\bar{q}V$ combinations. It can be written compactly as

$$\sigma_V = N_V \int_0^1 \int_0^1 \frac{dx_1 dx_2}{x_1 x_2} \delta(1 - \frac{\tau}{x_1 x_2}) P_{q\bar{q}V}(x_1, x_2),$$  

(4.3.4)

where $\tau$ has been defined in (4.2.6), the $\delta$-functions in (4.3.1) and (4.3.2) have been re-written using relation (4.2.5) and the overall normalization factor $N_V$ is given by

$$N_{W^\pm} = \frac{\pi^2 \alpha}{3 \sin^2 \theta_W} \quad \text{and} \quad N_{Z^0} = \frac{\pi^2 \alpha}{12 \sin^2 \theta_W \cos^2 \theta_W}.$$  

(4.3.5)

Since the Cabibbo factor and $n_q$ depend on the type of quarks in a specific subprocess, they have been absorbed in the joint $q\bar{q}$ probabilities, which are given by

$$P_{q\bar{q}W^\pm}(x_1, x_2) = \left( [u_1 \bar{d}_2 + c_1 \bar{s}_2] \cos^2 \theta_C + [u_1 \bar{s}_2 + c_1 \bar{d}_2] \sin^2 \theta_C \right) + (1 \leftrightarrow 2),$$  

(4.3.6)

similarly for $W^-$ with quark and antiquark densities interchanged, and

$$P_{q\bar{q}Z^0}(x_1, x_2) = \sum_f n_f q_1^f q_2^f + (1 \leftrightarrow 2).$$  

(4.3.7)

In the last two equations, the explicit $x$-dependence of the quark densities was left out ($q_i$ should be read as $q(x_i)$) and we assume contributions from the four lightest quarks only.

### 4.3.2 Higher Order QCD Corrections

In QCD, expression (4.3.4) merely represents the lowest order term of a perturbative expansion of the cross section in the strong coupling $\alpha_s$. Higher order corrections are introduced e.g. by graphs in which either the quark or the antiquark radiates a gluon before interacting, yielding the subprocess: $q + \bar{q} \rightarrow V + g$. In addition to $q\bar{q}$ pairs, one can also have other parton combinations in the initial state. An example is: $q + g \rightarrow V + q$, in which a virtual gluon in one of the hadrons splits into a $q\bar{q}$ pair of which the antiquark interacts with a quark from the other hadron.
4. \( W^\pm \) and \( Z^0 \) Production in \( p\bar{p} \) Collisions

In total, there are four types of parton subprocesses which contribute to the production of weak IVBs in \( p\bar{p} \) collisions, namely

\[
\begin{align*}
q + \bar{q} &\rightarrow V + \text{"partons"}, \quad (4.3.8a) \\
q(\bar{q}) + g &\rightarrow V + \text{"partons"}, \quad (4.3.8b) \\
q(\bar{q}) + q(\bar{q}) &\rightarrow V + \text{"partons"}, \quad (4.3.8c) \\
g + g &\rightarrow V + \text{"partons"}, \quad (4.3.8d)
\end{align*}
\]

where ‘partons’ stands short for quarks and gluons. The only subprocess of zeroth order in \( \alpha_S \) is the one discussed in the previous subsection, which is the lowest order term in (4.3.8a) with no partons in the final state. The lowest order term of the (anti)quark-gluon scattering subprocess in equation (4.3.8b), also referred to as the Compton subprocess, is of order \( \alpha_S \). Finally, the (anti)quark-(anti)quark and gluon-gluon scattering subprocesses in equations (4.3.8c) and (4.3.8d) are both of order \( \alpha_S^2 \).

Diagrams giving \( O(\alpha_S^2) \) corrections to the Born term are shown in figure 4.3 (virtual gluon graphs contribute via interference with the Born graph). The total cross section corrected to \( O(\alpha_S) \) was calculated in 1979, and is given by

\[
\sigma_V = N_V \int_0^1 \frac{dx_1}{x_1} \int_0^1 \frac{dx_2}{x_2} \left\{ P_{q\bar{q}V}(x_1,x_2,\mu) \left[ \delta(1 - \hat{t}) + \frac{2}{3\pi} \alpha_S(\mu) \theta(1 - \hat{t}) 2 f_q(\hat{t}) \right] \\
+ P_{g\bar{q}V+g\bar{q}V}(x_1,x_2,\mu) \frac{1}{4\pi} \alpha_S(\mu) \theta(1 - \hat{t}) f_g(\hat{t}) \right\}, \quad (4.3.9)
\]

where \( \hat{t} \equiv \tau/(x_1x_2) \) and the functions \( f \) are defined as

\[
\begin{align*}
f_q(z) &= \left\{ \frac{3}{2} \frac{1}{1 - z} \frac{1}{(1 - z)} \frac{\ln(1 - z)}{1 - z} - 3 - 2z + \left( \frac{1}{2} + \frac{2}{3\pi} \right) \delta(1 - z) \right\} \\
f_g(z) &= \left\{ \frac{9}{2} z^2 - 5z + \frac{3}{2} + \left( z^2 + (1 - z)^2 \right) \ln(1 - z) \right\}.
\end{align*}
\]

The joint \( (q\bar{q} + \bar{q}q) \) probabilities are given by

\[
P_{q\bar{q}w^+ + \bar{q}q}w^+(x_1,x_2,\mu) = \left( u_1 + c_1 + \bar{d}_1 + \bar{s}_1 \right) g_2 + (1 \leftrightarrow 2), \quad (4.3.11)
\]
similarly for \( W^- \) with quark and antiquark densities interchanged, and

\[
P_{q\bar{q}z^0 + \bar{q}q}z^0(x_1,x_2,\mu) = \sum_f n_f \left( q_f^1 + \bar{q}_f^1 \right) g_2 + (1 \leftrightarrow 2). \quad (4.3.12)
\]

In equation (4.3.9), the term containing the \( \delta \)-function is the Born term. The other ones correspond to the creation of an additional (anti)quark or gluon in the final state. They require the CM energy squared of the two initial state partons to be larger than the IVB mass squared: \( \hat{s} \geq M^2 \). This condition is generated by the step-functions, \( \theta(1 - \hat{t}) \).
4.3 The Total $W^\pm$ and $Z^0$ Production Cross Sections

\[
\begin{array}{c}
\text{a) } \\
\text{b) } \\
\text{c) }
\end{array}
\]

Figure 4.3: Diagrams giving $O(\alpha_S)$ corrections to the Born term. Real-gluon corrections from: a) the annihilation and b) the Compton subprocesses, and virtual-gluon corrections (c).

The parton densities now depend on the four-momentum transferred between the scattering partons, $Q^2$. This dependence has been indicated by a scale $\mu$, known as the factorization scale ($q_f$ should be read as $q(x_f, \mu)$). In QCD, the factorization theorem\textsuperscript{68} states that the $Q^2$-dependent parton densities are process independent. This theorem is of great importance, because it allows to employ the results on quark and gluon densities measured by deep inelastic lepton production experiments and make actual predictions for the $W^\pm$ and $Z^0$ production cross sections using relation (4.3.9). The parton densities, which are usually measured at low $Q^2$ values ($\lesssim 10$ GeV$^2$), have to be evolved by the Altarelli-Parisi equations to $Q^2$ values relevant for weak IVB production ($\sim M^2$).

The renormalization scale, which appears in the running coupling $\alpha_S$, can be chosen with some freedom around the natural physical scale of the process\textsuperscript{56} giving variations in the cross section of $O(\alpha_S^2)$. For simplicity, the factorization and renormalization scale have been chosen to be equal. In the case of IVB production, the scale is usually chosen as: $\mu = M^2$, where $M$ is the vector boson mass. However, the transverse momentum ($p_T$) distribution of weak IVBs with $<p_T^2> \ll M^2$ would perhaps suggest a lower scale, hence a larger $\alpha_S$.

Recently, the lowest order terms ($O(\alpha_S^2)$) of the subprocesses (4.3.8) (c)$\textsuperscript{69}$ and (d)$\textsuperscript{70}$ and the most important part of the $O(\alpha_S^2)$ term in (4.3.8a)$\textsuperscript{71}$ have been computed. The missing parts of $O(\alpha_S^2)$ are expected to be small at SppS energies$\textsuperscript{70}$. 

4. W± and Z⁰ Production in pp Collisions

Table 4.1: Total cross sections for W and Z production (σW and σZ) at the CERN Sp⁰S and the Tevatron at Fermilab, and the ratio Rσ = σW/σZ. The results are given to 0th, 1st and 2nd order in αs and have been computed with two different sets of structure functions.

<table>
<thead>
<tr>
<th>⟨s⟩ (GeV)</th>
<th>str. function</th>
<th>O(αs⁰)</th>
<th>O(αs¹)</th>
<th>O(αs²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>630</td>
<td>MTS</td>
<td>σW (nb)</td>
<td>4.77</td>
<td>6.13</td>
</tr>
<tr>
<td></td>
<td></td>
<td>σZ (nb)</td>
<td>1.50</td>
<td>1.93</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Rσ</td>
<td>3.18</td>
<td>3.18</td>
</tr>
<tr>
<td>630</td>
<td>DFLM260</td>
<td>σW (nb)</td>
<td>4.43</td>
<td>5.78</td>
</tr>
<tr>
<td></td>
<td></td>
<td>σZ (nb)</td>
<td>1.37</td>
<td>1.79</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Rσ</td>
<td>3.23</td>
<td>3.23</td>
</tr>
<tr>
<td>1800</td>
<td>MTS</td>
<td>σW (nb)</td>
<td>15.26</td>
<td>18.70</td>
</tr>
<tr>
<td></td>
<td></td>
<td>σZ (nb)</td>
<td>4.61</td>
<td>5.67</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Rσ</td>
<td>3.31</td>
<td>3.30</td>
</tr>
<tr>
<td>1800</td>
<td>DFLM260</td>
<td>σW (nb)</td>
<td>15.89</td>
<td>19.73</td>
</tr>
<tr>
<td></td>
<td></td>
<td>σZ (nb)</td>
<td>4.81</td>
<td>5.98</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Rσ</td>
<td>3.30</td>
<td>3.30</td>
</tr>
</tbody>
</table>

4.3.3 Numerical Predictions

We now consider numerical predictions for the total cross sections for W± and Z⁰ production in pp collisions, and their ratio. For the complete calculation including the O(αs⁰) and all available O(αs¹) corrections, the program ZWPROD⁷² has been used. It computes matrix elements in the MS-scheme and uses for the running coupling αs the one defined to two loop order. Mass factorization has been performed in the DIS-scheme, requiring the use of parametrizations of parton distributions in the same scheme⁷³. The calculation of the subprocess cross section in next-to-leading order (NLO) QCD motivates the use of parton distribution functions evolved also in NLO, of which a few have become available recently. The ones employing the DIS-scheme are the DFLM⁷⁴ and MT⁷⁵ distributions. In the calculation, we have used the following input parameters: MZ = 91.172 GeV, MW = 80.00 GeV, sin²θW = 0.230 and sin²θC = 0.05.

In table 4.1 are listed the W± and Z⁰ total cross sections and the ratio Rσ in 0th, 1st and 2nd order αs. For the parton densities, we have used the parametrizations of MT (set S) with a QCD scale parameter Λ(2nd order, 4 light quark flavours) = 212 MeV, and of DFLM with Λ(2,4) = 260 MeV. At ⟨s⟩ = 630 GeV (1800 GeV) the results show a 30% (24%) and 9% (10%) increase of the cross section from a 1st and 2nd order correction respectively. The smaller NLO correction with respect to the LO correction confirms the perturbative character of the cross section expansion. Note that Rσ is practically insensitive to higher order corrections. The origin to this phenomenon is explained in reference 71, where the authors remark that the first and second order corrections are dominated by the δ(1 − t) term appearing in the q̅q subprocess (see also equation (4.3.9) and the definition of the f_q-
4.3 The Total $W^\pm$ and $Z^0$ Production Cross Sections

Table 4.2: The ratio: $\sigma_Z(\mu=(xM)^2)/\sigma_Z(\mu=M^2)$, displaying the scale dependence of the 0th, 1st and 2nd order cross section.

<table>
<thead>
<tr>
<th>$\sqrt{s}$ (GeV)</th>
<th>$x$</th>
<th>$O(\alpha_S^0)$</th>
<th>$O(\alpha_S^1)$</th>
<th>$O(\alpha_S^2)$</th>
</tr>
</thead>
<tbody>
<tr>
<td>630</td>
<td>1/2</td>
<td>1.07</td>
<td>1.04</td>
<td>1.01</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>0.94</td>
<td>0.97</td>
<td>0.99</td>
</tr>
<tr>
<td>1800</td>
<td>1/2</td>
<td>0.98</td>
<td>1.01</td>
<td>0.99</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>1.01</td>
<td>1.00</td>
<td>1.01</td>
</tr>
</tbody>
</table>

function in (4.3.10)). Because this contribution is just an overall normalization factor for both $\sigma_W$ and $\sigma_Z$, it cancels in the ratio.

There are several sources of uncertainty on these predictions. One already mentioned earlier, is associated with the renormalization and factorization scale. There is no clear consensus among theorists on what a proper choice for the scale would be. However, we do know that changing the scale in the cross section corrected up to $O(\alpha_S^3)$ can only lead to variations in the result of $O(\alpha_S^2)$ (we have neglected the fact that the 2nd order calculation is not entirely complete). Extrapolating the trend of a 30% 1st order correction and a 10% 2nd order correction, we expect variations in the cross section due to a change in scale not to exceed a few percent. This is confirmed by the results listed in table 3.2 computed with the MT structure functions, which shows the ratio of $\sigma_Z$ evaluated at the scale $\mu = (xM)^2$ and $\sigma_Z$ evaluated at $\mu = M^2$ for $x = 1/2$ and $x = 2$. Since the numerical values are similar for the ratio of $W$ cross sections, the scale dependence cancels in $R_S$. We observe a $\pm 1\%$ change in the $O(\alpha_S^2)$ result upon varying the scale within the quoted range of $x$. In reference 70, the effect of varying the factorization and renormalization scales independently was investigated, which showed in the extreme case a 2% (4%) change in the $O(\alpha_S^2)$ result at $\sqrt{s} = 630$ GeV (1800 GeV). We take these last numbers as a theoretical uncertainty due to the choice of scale, because on the one hand the $\mu$-interval implied by our choice of $x$ is arbitrary (and may not be wide enough) but on the other hand the extreme choice for the factorization and renormalization scale is probably overly pessimistic.

Another source of uncertainty is introduced by the choice of structure functions. The results in table 4.1 show a $\sim 6\%$ change in the cross sections when switching from MT to DFLM parton distributions, which is almost independent of the order of the calculation. In table 4.3, this point is illustrated with a few more sets of structure functions, where cross sections have been computed in lowest order only. It is a common practice, however incorrect, to apply a variety of parton distributions to a given physical process and then cite the range of results obtained as the ‘theoretical error’, because: (i) some distributions are known to disagree with current data, and in many cases deviate from the correct result in the same direction rather than ‘bracketing’ the right answer, (ii) the treatment of experimental errors and corrections often differs considerably for the various parton distributions (apart from the MT distributions, none of them takes into account the experimental systematic errors)\textsuperscript{76}. Furthermore, the precision of NLO calculations would be spoiled by the application of parton densities evolved in LO only.
4. $W^\pm$ and $Z^0$ Production in $p\bar{p}$ Collisions

Table 4.3: Zeroth order cross sections and their ratio for different sets of structure functions.

<table>
<thead>
<tr>
<th>str. function</th>
<th>$\sqrt{s} = 630$ GeV</th>
<th>$\sqrt{s} = 1800$ GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>$\sigma_W^0$ (nb)</td>
<td>$\sigma_Z^0$ (nb)</td>
</tr>
<tr>
<td>MTS</td>
<td>4.77</td>
<td>1.50</td>
</tr>
<tr>
<td>DFLM260</td>
<td>4.43</td>
<td>1.37</td>
</tr>
<tr>
<td>DO$^{177}$</td>
<td>4.69</td>
<td>1.36</td>
</tr>
<tr>
<td>GHR$^{78}$</td>
<td>4.39</td>
<td>1.37</td>
</tr>
<tr>
<td>HMRSB$^{79}$</td>
<td>4.95</td>
<td>1.51</td>
</tr>
<tr>
<td>HMRSE$^{79}$</td>
<td>5.31</td>
<td>1.59</td>
</tr>
</tbody>
</table>

With these remarks in mind, the only results useful to us are the ones obtained with the MT and DFLM structure functions. The authors of the MT distributions made the most robust analysis at present (accounting for systematic errors in the various data sets and for kinematical cuts applied by the different experiments) using the largest statistics available. Therefore, instead of taking the average of the MT and DFLM results, we prefer to use the cross sections obtained with the MT distributions as central values assuming a 5% 'theoretical uncertainty' due to structure functions.

For $R_\sigma$ the situation is different, because it is hardly affected by higher order corrections, so that in principle one could apply other parton densities with LO evolution only. In chapter eight, we will discuss a new criterion based on measurements of the structure function ratio, $F_2^L/F_2^R$, which can be used to constrain our choice of parametrizations of parton densities for the evaluation of $R_\sigma$.

In chapter three, we presented the following uncertainties on the value of $M_W$ and $\sin^2\theta_W$: $\Delta M_W = \pm 0.21$ GeV and $\Delta \sin^2\theta_W = 0.002$. The former leads to a variation in $\sigma_W$ of: $\delta \sigma_W = \pm 18$ pb (38 pb) at $\sqrt{s} = 630$ GeV (1800 GeV), while the latter creates a change in $\sigma_Z$ of: $\delta \sigma_Z = \pm 5$ pb (15 pb) at $\sqrt{s} = 630$ GeV (1800 GeV). In both cases at $\sqrt{s} = 630$ GeV and 1800 GeV the change in $R_\sigma$ is: $\Delta R_\sigma = 0.01$.

In conclusion, we use the following numerical predictions for the $W^\pm$ and $Z^0$ total cross sections

$$
\begin{align*}
\sigma_W(630) &= 6.67 \pm 0.13 \text{ (scale)} \pm 0.33 \text{ (str.functions)} \pm 0.02 \text{ (} M_W \text{) nb}, \\
\sigma_Z(630) &= 2.09 \pm 0.04 \text{ (scale)} \pm 0.10 \text{ (str.functions)} \pm 0.01 \text{ (} \sin^2\theta_W \text{) nb}, \\
\sigma_W(1800) &= 20.53 \pm 0.82 \text{ (scale)} \pm 1.03 \text{ (str.functions)} \pm 0.04 \text{ (} M_W \text{) nb}, \\
\sigma_Z(1800) &= 6.20 \pm 0.25 \text{ (scale)} \pm 0.31 \text{ (str.functions)} \pm 0.02 \text{ (} \sin^2\theta_W \text{) nb}.
\end{align*}
$$

(4.3.13)

These values correspond to $R_\sigma(630) = 3.19$ and $R_\sigma(1800) = 3.31$. At this stage, we refrain from quoting uncertainties on $R_\sigma$. These will be discussed in chapter eight.
5 The UA1 Apparatus
5. The UA1 Apparatus
5.1 Introduction

In the previous two chapters, we discussed theoretical input required for the determination of $\Gamma_{W}^{tot}$, and SM predictions concerning vector boson production and decay. The coming three chapters are dedicated to a discussion of UA1 data taken in the 1988 and 1989 runs. In the present one, a brief overview of the CERN collider is given followed by a description of the UA1 detector and a short discussion of the data acquisition and the luminosity measurement.

5.2 The CERN Proton-Antiproton Collider

5.2.1 Luminosity

In the introductory chapter, we mentioned what motivated physicists to convert the SPS at CERN into a p$\bar{p}$ storage ring, namely: the direct production of $W^{\pm}$ and $Z^{0}$ bosons in p$\bar{p}$ collisions. Among many problems, a major obstacle to the p$\bar{p}$ collider concept was the difficulty to produce the required amount of interactions to detect the weak IVBs. The total number of produced vector bosons, $N_{V}$, in a certain amount of running time at a collider is proportional to the cross section for the production process $\sigma_{V}$, and to the integrated luminosity of the machine $\int \mathcal{L} \, dt$

$$N_{V} = \sigma_{V} \times \int \mathcal{L} \, dt. \quad (5.2.1)$$

Since the cross sections were known to be of the order of a few nanobarns, an integrated luminosity of at least 100 nb$^{-1}$ was needed to produce vector bosons for detection and even more if one wanted to measure their properties. To collect 100 nb$^{-1}$ or more in a reasonable amount of running time, an instantaneous luminosity of about $\mathcal{L} = 10^{29} - 10^{30}$ cm$^{-2}$s$^{-1}$ is required.

The luminosity depends on the number of protons and antiprotons per bunch $N_{p}$ and $N_{\bar{p}}$, the number of bunches $n$ in either beam, the effective cross sectional area of the beams $A$, and the revolution frequency $f$, as follows

$$\mathcal{L} = fn \frac{N_{p}N_{\bar{p}}}{A}. \quad (5.2.2)$$

If the beams collide head on, $A$ is given by

$$A = 4\pi \sigma_{h} \sigma_{v}. \quad (5.2.3)$$

The quantities $\sigma_{h}$ and $\sigma_{v}$ are the rms horizontal and vertical beam dimensions. Therefore, in order to obtain a high luminosity, the beam dimensions at the interaction point should be small and the number of protons and antiprotons should be large. The former objective is fulfilled by using special focusing magnets near the interaction region. At UA1, the beams are squeezed to lateral dimensions less than 1 mm. The revolution time for a proton of 270 GeV in the SPS is 23 $\mu$s corresponding to a revolution frequency of: $f \approx 43$ kHz. The number of bunches per beam was foreseen to be six, however, during the first years of operation the actual number used, was three. Only in 1988, the switch from three to six bunches per beam was made.
With this set of parameters, one needs \( N_p N_{\bar{p}} = 10^{22} \) in order to reach the high luminosities quoted above. Since unequal proton and antiproton bunch densities lead to beam instabilities and hence to a reduced lifetime of the beams, it is desirable to have \( N_p = N_{\bar{p}} \). It is relatively easy to obtain \( 10^{11} \) protons per bunch. However, to reach these bunch densities for antiprotons, special techniques have to be applied. This problem will be addressed in the next subsection.

5.2.2 Accumulation of Antiprotons

Since antiprotons have to be created in high energy collisions, it is rather difficult to produce an antiproton beam of high intensity. At CERN they are obtained by dumping a dense 26 GeV proton beam coming from the Proton Synchrotron (PS) onto a target. From each burst of about \( 10^{13} \) protons, \( 5 \times 10^6 \) antiprotons are obtained, which have been created with a wide range of angles and an energy spread which could never be accommodated within the acceptance of a machine such as the SPS. They are therefore guided through a beam transport system to a large aperture storage ring, the Antiproton Accumulator (AA), where their motion is ‘cooled’. Cooling allows the beam dimensions to be reduced or the phase-space density to be increased by many orders of magnitude.

Several methods of beam cooling exist. The one applied at CERN is called ‘stochastic cooling’. It basically consists of a beam pick-up, which detects random fluctuations in the centre of gravity of a sample of beam circulating in a storage ring. The signal is amplified and sent to a kicker magnet on the opposite side of the ring, where it arrives at the same time as the sample and corrects the error. S. van der Meer\(^8^5\) was the first to show that continuously correcting the mean of a sample slowly reduces the beam size. The antiprotons which are injected into the AA, are pre-cooled in the outer part of the wide AA vacuum chamber for about 2 seconds. After precooling, they are stacked in the inner part of the ring, where cooling continues. With a stacking rate of \( 5 \times 10^9 \) \( \bar{p} \)/hr, up to \( 10^{11} \) antiprotons can be collected per day.

In order to remain competitive with the Tevatron at Fermilab in the U.S., which came online in 1987, a second ring, the Antiproton Collector (ACOL), was added to the AA. With ACOL, the number of antiprotons that can be captured from the production target has substantially increased, boosting the stacking rate by an order of magnitude.

5.2.3 Beam Transfer and Acceleration

Figure 5.1 shows that part of the CERN accelerator complex which is relevant to the \( p\bar{p} \) collider. It illustrates the transfer of the beams and the different stages of acceleration. The CERN 50 MeV linear accelerator, the 1 GeV PS Booster, and the 26 GeV PS, together produce about \( 10^{13} \) protons per 2.4 seconds. In the first stage, these protons are used to produce antiprotons as described above.

After antiprotons have been stacked for many hours, the PS and SPS prepare for a fill. Subsequently, three (six) proton bunches are sent to the SPS at 2.4 s intervals and stored at 26 GeV. These bunches are placed azimuthally symmetric around the SPS circumference. Then a single antiproton bunch is unstacked from the AA and injected into the PS at 3.5 GeV, accelerated to 26 GeV, and injected into the SPS on the correct
Figure 5.1: Layout of the CERN accelerator complex. The fat lines indicate the modifications made to the existing complex in order to make pp collisions possible. Also shown are the different stages of acceleration of the protons and antiprotons.

azimuthal position to within a fraction of a nanosecond for collision in the detectors. This is followed at 2.4 s intervals by two (five) more antiproton bunches. The machine then accelerates both the proton and antiproton beams to storage energy, and keeps them in storage for up to 24 hours whilst the next batch of antiprotons is being prepared in the AA.

Table 5.1 shows the beam energies and integrated luminosities delivered by the machine for the years 1982—1989. Notice the strongly increased integrated luminosity in the last two years thanks to the introduction of ACOL. In 1987, a technical run was made to test the upgraded version of the machine.

<table>
<thead>
<tr>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>Beam energy (GeV)</td>
<td>273</td>
<td>273</td>
<td>315</td>
<td>315</td>
<td>315</td>
<td>315</td>
</tr>
<tr>
<td>Number of bunches per beam</td>
<td>3</td>
<td>3</td>
<td>3</td>
<td>3</td>
<td>6</td>
<td>6</td>
</tr>
<tr>
<td>Int. luminosity delivered (nb⁻¹)</td>
<td>28</td>
<td>153</td>
<td>395</td>
<td>655</td>
<td>3372</td>
<td>4760</td>
</tr>
</tbody>
</table>
5. The UA1 Apparatus

5.3 The UA1 Detector

5.3.1 An Overview

In 1978, a series of proposals for detectors at the p\bar{p} collider was approved, including the one described here, named UA1 for Underground Area 186. It is located 20 meters underground in one of the Long Straight Sections, LSS5, of the SPS ring.

The UA1 detector is a general purpose detector designed to study p\bar{p} collisions with a special emphasis on the detection of weak IVBs through their leptonic decay mode. It provides tracking and momentum measurements of charged particles in the Central Drift chamber (CD), hadronic calorimetry, external drift chambers for muon identification and tracking, and iron absorber walls, instrumented with limited streamer tubes for improved muon measurements. In the data taking period up to 1985, electromagnetic calorimetry was present as well.

A side view of the apparatus along the beam direction is shown in figure 5.3 on the opposite page. It refers to the configuration as in 1985. Protons and antiprotons in bunches of about 30 cm collide at the centre of the apparatus. The interaction region is inside a stainless steel beam pipe of ~130 mm inner diameter with 0.15 mm thick walls. The pipe is corrugated with a maximum outer diameter of 146 mm. The interaction region is surrounded by successive layers of special purpose detectors. The tracking and magnetic analysis is made in the CD, where charged particle tracks are detected in the pseudo-rapidity range \(|\eta| \leq 3\). The electromagnetic calorimeters were absent in the 1988 and 1989 runs, leaving empty space. Outside this region is a 13 cm aluminium coil of a dipole magnet, which produces a uniform horizontal field of 0.7 T in the region of the CD. Outside the coil is a large iron-scintillator sandwich hadron calorimeter, which also serves as the return yoke of the magnet. This is in turn surrounded by iron absorber drift chambers used to identify muons that pass through the hadron calorimeters and iron absorber. The forward calorimeters and very forward calorimeters were absent during the 1988 and 1989 runs. The iron absorber in the side and forward regions has been instrumented with limited streamer tubes to aid in the muon momentum measurement. Before describing the CD, hadron calorimeters and muon chambers in more detail, we first introduce the coordinate system used in UA1 shown below in figure 5.2.

\[
\begin{align*}
    p_T &= \sqrt{(p_y^2 + p_Z^2)}, \\
    \varphi &= \text{atan} \left( \frac{p_y}{p_Z} \right), \\
    \theta &= \text{atan} \left( \frac{p_T}{p_x} \right), \\
    \lambda &= \text{atan} \left( \frac{p_Z}{\sqrt{(p_x^2 + p_y^2)}} \right), \\
    y &= \frac{1}{2} \ln \left( \frac{E + p_x}{E - p_x} \right), \\
    \eta &= -\ln \tan \left( \frac{\theta}{2} \right), \\
    \Delta R &= \sqrt{\Delta \eta^2 + \Delta \varphi^2}
\end{align*}
\]

Figure 5.2: Definition of coordinates in UA1. The x-axis is the direction of the antiproton beam; the z-axis is the direction of the magnetic field; \(\varphi\) is the azimuthal angle around the beam axis; \(\theta\) is the polar angle with respect to the outgoing antiproton beam; \(\lambda\) is the dip angle with respect to the plane of curvature (xy-plane); \(y\) is the rapidity; \(\eta\) is the pseudorapidity; \(\Delta R\) is the distance in the \(\eta\varphi\)-plane.
Figure 5.3: A side view of the UA1 apparatus.
5. The UA1 Apparatus

5.3.2 The Central Detector

The CD forms the core of the charged particle identification in the UA1 experiment. It is composed of six independent semi-cylindrical chambers assembled to form a cylinder, 5.8 m long and 2.3 m in diameter, covering the polar angle from 5° to 175° with respect to the beam direction (figure 5.4). Each semi-cylinder is separated into drift volumes by wire planes. The 6125 sense wires and 17000 field shaping wires are parallel to the magnetic field. Cathode planes (typical voltage ~30 kV) alternate with anode planes. Figure 5.5 shows how the sense wires (at ~1.5 kV) and field shaping wires (at ~2.5 kV) in the anode planes are combined in three layers. This arrangement ensures that each wire plane collects signals from only one of the neighbouring drift volumes.

The chamber is filled with a gas mixture of 40% argon and 60% ethane, which is easily ionizable. If charged particles traverse the drift volume, they liberate electrons from the gas molecules, which will drift in the uniform electric field towards the sense wires. Near a sense wire the electric field gradient increases significantly, which allows the drift electrons to gain sufficient energy to ionize the gas and to form an avalanche. This effect, called gas amplification, is controlled by the field shaping wires. The charge collected on a sense wire produces a pulse, which can be read out at both ends of the wire.

The drift field is set to 1.5 kV/cm, giving a saturated drift velocity of 53 mm/μs. A Lorenz force, induced by the electric field and the magnetic field, rotates the drift direction of the electrons by an angle of about 23°. The spacing of the drift volumes has been designed to keep the maximum drift time just below the bunch crossing time of 3.8 μs. With the drift velocity given above, this corresponds to a wire gap of 18 cm, which allows for a drift angle of 23°. The arrangement of the wire planes minimizes the regions in which tracks coming from the vertex go along the drift direction. This is illustrated in figure 5.6.
The ends of the sense wires are connected to preamplifiers, which are mounted on the outside surface of the CD. The signals coming from the preamplifiers are sent to the counting room, where they are digitized. A readout system records continuously the information on the drift time in bins of 4 ns, the pulse height in bins of 32 ns and the charge division along a wire. From drift times and wire positions, the track coordinates in the plane perpendicular to the wires (xy-plane) are determined. The coordinate along the wire is measured by comparing the charge collected on either end of the wire. The position resolution in the xy-plane is of the order of 300 \( \mu \text{m} \). The resolution obtained in the z-direction is of a much lower quality and is of the order of several centimetres.

One event contains roughly 10000 digitizings, which makes pattern recognition in the CD a complicated task. A special procedure has been developed\(^8\), which uses a parabolic path to build a chain of consecutive points. The basic idea relies on the observation that neighbouring points belonging to the same track are in general closer than points on nearby tracks. A typical momentum resolution of

\[
\frac{\Delta p}{p^2} \approx 0.015 \text{ GeV}^{-1}
\]  
(5.3.1)
is obtained. However, the resolution of a track depends strongly on its position in the CD.

A final word about some changes made to the CD in preparation of the 1988 and 1989 runs. The instantaneous luminosity during these runs was about a factor 6 higher than the one in the runs before. Under unchanged circumstances the currents drawn by the CD would surpass a safety limit. Therefore, the gas amplification was lowered by a factor four, which was partially compensated by an increase in electronic gain of a factor three. The resolutions quoted above correspond to the situation as in 1988 and 1989.

5.3.3 The Hadron Calorimeter

The magnetic curvature analysis of charged tracks in the CD is complemented by the measurement of both charged and neutral particle energies in the calorimeter. In particular, for events with large particle densities, energy flow measurements remain significant whilst curvature determinations become difficult. In general, a distinction is made between electromagnetic and hadron calorimeters. The former type is used for the detection of particles with an early shower development, such as electrons and photons, while the latter detects particles with a late shower development. Originally, the UA1 detector contained both types of calorimetry. In 1985 however, it was decided to replace the electromagnetic calorimeter by an upgraded version, which was to be installed before the 1988 run. Due to technical problems the installation was delayed and finally cancelled. Therefore, during the 1988 and 1989 runs the UA1 detector contained a hadron calorimeter only.
The hadron calorimeter\textsuperscript{89,90} consists of instrumented iron, which serves as a return yoke for the dipole magnet (figure 5.7). It covers the polar angle range: $5^\circ < \theta < 175^\circ$ (or correspondingly, to a pseudorapidity range: $-3 < \eta < +3$) and is constructed from two different types of modules: 16 C-modules covering angles from 25$^\circ$ to 155$^\circ$ and 12 I-modules forming the end-caps which extend the angular coverage down to 5$^\circ$. Each module consists of several stacks of iron and scintillator plates. Light collected in the scintillator plates is read out on both sides of a stack via BBQ wavelength shifter bars and transported through light guides to photomultipliers.

Eight C-modules on each side of the experiment form the barrel part of the calorimeter. They are segmented into 12 azimuthal sections, each with two samplings in depth corresponding to 5 interaction lengths in total. Figure 5.8 shows one C-module with the light guides and photomultipliers for two cells as indicated. The empty space between the CD and the hadron calorimeter was previously occupied by the electromagnetic calorimeter.

The 12 I-modules make up part of the end-caps of the detector. Each end, consisting of six vertical sections, is segmented into 36 rectangular stacks, with further subdivision into four smaller stacks for the 16 stacks nearest to the beam (figure 5.7 (b)). The small stacks (small I's) are read out by only one BBQ bar, whereas the large stacks (large I's) have two-sided readout as for the C-modules. Each I-cell contains two samplings in depth corresponding to 7 interaction lengths in total.

The energy resolution was measured using prototype modules exposed to hadron beams in the momentum range 0.7 to 90 GeV. The dependence of the resolution on the particle energy can be parametrized as\textsuperscript{90}

$$\frac{\sigma(E)}{E} = \frac{0.80}{\sqrt{E}}, \quad (5.3.2)$$

with $E$ in units of GeV.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure5.7.png}
\caption{(a) Top view of the Barrel (C-modules) and End-Cap (I-modules) hadron calorimeter, and (b) the granular structure of the I-modules.}
\end{figure}
Figure 5.8: View of a C-module showing the exterior of the CD, the magnet coil and some features of the stack readout.

5.3.4 The Muon Detection System

The outer shell of the apparatus consists of a large drift chamber array, which constitutes the muon detection system (particles coming from the interaction point, other than muons, are usually absorbed in the hadron calorimeter or the extra iron absorber behind the calorimeter). There are in total 34 chambers which cover about 70% of the solid angle. Each chamber, which measures $4 \times 6 \text{ m}^2$, contains four layers of drift tubes, two by two in orthogonal projections. Two adjacent planes are staggered in order to solve left-right ambiguities.
With the exception of ten bottom chambers, the muon chambers are grouped in pairs which are separated by 60 cm to form a muon chamber module. This separation provides a lever arm which allows to measure not only the position, but also the direction of a track. The organization of drift tubes in a muon chamber and muon chambers in a muon chamber module is illustrated in figure 5.9.

Due to space limitations, the modules in the bottom region of the detector consist of only one chamber. They differ in size from the other chambers in order to fit between the rails which support the detector. There are two large chambers of $5 \times 6 \text{ m}^2$ and eight smaller ones of $1 \times 6 \text{ m}^2$.

![Diagram of muon drift chamber](image)

**Figure 5.9:** Schematic view of a muon drift chamber module.

The drift tubes are made of extruded aluminum and measure $45 \times 150 \text{ mm}^2$ with a variable length. Each tube contains a $50 \mu\text{m}$ stainless steel wire in the centre. A cross sectional view of a single drift tube is given in figure 5.10. The particular shape of the tube around the anode wire restricts the sensitive drift volume to a region of equal length field lines. Field shaping cathodes at the edges of the tube constrain the electric field such, that a constant drift velocity of $52 \text{ mm/}\mu\text{s}$ is maintained over the whole width of the tube. The maximum drift time, which is recorded in 8 ns bins, is $1.4 \mu\text{s}$. 
5.4 Trigger and Data Acquisition

5.4.1 Trigger Levels and Readout Structure

With six bunches each of protons and antiprotons the CERN Collider provides bunch crossings in the UA1 detector every 3.8 μs. Although not every bunch crossing yields an interaction, the high inelastic pp cross section, of which about 34 mb is seen in the UA1 apparatus, produces an interaction rate of about 70 kHz at 2 × 10^{30} cm^{-2}s^{-1}. The event rate for processes of interest, however, is much lower. For example, the cross section for Z production times branching ratio for a Z decay into two muons is of the order of 50 pb. If we neglect a further detection efficiency, then at the same luminosity as above this process occurs at a rate of 10^{-4} Hz or only once every three hours! Clearly, a sophisticated event filtering procedure is needed. This is accomplished by the use of a multilevel trigger system which pipelines the trigger decision (accept/reject event) from fast unintelligent low level triggers to more sophisticated higher level triggers with longer execution times.

Three levels of triggering are used in UA1. A schematic representation of the readout structure and the different trigger levels is given in figure 5.11. Event rates at different stages in the readout have been indicated on the right of the diagram for a luminosity of 1.5 × 10^{30} cm^{-2}s^{-1}, typical for the 1988 and 1989 runs. The main bottle neck in the system is the enormous volume of data produced by the CD. Digitization takes less than 3.8 μs, but data reduction and readout require 30 ms. Therefore, the first and second level triggers must reduce the trigger rate to well below 30 Hz and cannot use the CD information.
Two separate two-level trigger systems are used, one based on the muon chambers and the other on the calorimeter. Both systems operate on digitized information only. The combined process of digitization of detector signals, generation of a first level trigger decision and clearing of detector elements in case of a ‘reject event’ decision, takes less than 3.8 $\mu$s, so that no dead time (the time during which the detector is insensitive to the occurrence of new events) is created at level 1. The typical reduction in event rate achieved at this level is about a factor 1000. Upon a valid first level trigger, the readout of the CD is started while muon and calorimeter data are stored in a double buffer.

The second level trigger operates on this buffer with a mean execution time of 8 ms and typically reduces the rate by a factor 5. If an event is rejected by the second level trigger, further readout of the CD is inhibited and the buffer is cleared. Upon accept, the information from the different detector elements is read out in parallel and combined by an ‘event builder’ to construct a complete event. During readout the second level trigger can operate on one more event.

![Trigger and Data Acquisition Diagram](image)

Figure 5.11: Schematic representation of the readout structure and the different trigger levels in the UA1 experiment. Event rates at different stages in the process have been indicated for a luminosity that was typical for the 1988 and 1989 runs.
5. The UA1 Apparatus

The event builder sends the event to cassette (the so-called ‘normal’ data stream) and passes a copy to the third level trigger. The purpose of this trigger is twofold. First of all, it selects all events that are considered interesting while reducing the rate by a factor 20 to allow a quick data analysis. Secondly, it provides a possibility for offline monitoring (in addition to the usual online monitoring) of the proper functioning of the detector with a fast feedback. Events passing the third level trigger are referred to as the ‘special’ data stream. Under normal running conditions the deadtime generated by the entire system is about 12%. Special tests showed a threefold increase of the deadtime for runs in single buffer mode.

In the following, each of the trigger levels will be discussed in more detail. Usually several trigger conditions are active simultaneously. The ones most commonly used during the 1988 and 1989 runs are inclusive one- and di-muon triggers sometimes backed up by an additional calorimeter jet requirement.

5.4.2 The First Level Trigger

At the basis of the trigger system is a pretrigger, which consists of fast scintillation hodoscopes placed symmetrically around the nominal interaction point outside the end-cap calorimeters. A valid pretrigger condition consists of a coincidence between counter hits from both arms strobed by a machine crossing signal. Test runs on machine crossings have shown that this trigger introduces very little bias. Its efficiency is essentially 100%.

Clearly, with less than 3.8 μs available at level 1, the calorimeter trigger cannot handle the full granularity and segmentation of the calorimeter. By adding signals from groups of neighbouring cells with analog electronics, the 1246 calorimeter cells are mapped onto 144 trigger channels which are digitized with 8-bit flash ADCs. The digital information is subsequently sent to the digital trigger logic, where raw data words are converted to transverse energies by addressing lookup memories. Local and global energy sums are then compared to threshold values for jets and the total energy for the event respectively.

Since digitization of drift times measured in the muon chambers takes about 30 μs, this information is not available at level 1. Instead, the muon trigger employs the information on which drift tubes have been hit. The width of the tubes (15 cm) fixes the granularity of the pattern from which the muon trigger derives its decision. Track finding is performed in separate projections, as illustrated in figure 5.9. Signals from groups of ten wires in a projection are combined (hard wired) to form an address in a lookup table, which determines whether tube combinations correspond to a track coming from the centre of the detector (within ± 150 mrad of the nominal vertex). A valid track requires the reconstruction of both track projections in a module, with an exception for the bottom chambers where only one projection can be measured. The pointing requirement discriminates against hadronic punch-through and low momentum muon tracks that are more likely to be bent by the magnetic field and by multiple scattering.
5.4 Trigger and Data Acquisition

Figure 5.12: Angular coverage of the muon chambers for different trigger regions. The two regions mainly used in the 1988 and 1989 runs are the barrel (solid) and F3 (dotted) areas.

All calorimeter backstacks (the cells closest to the muon chambers) have been equipped with discriminators which fire when the deposited energy exceeds a certain noise threshold. A muon candidate is accepted by the first level trigger, if a valid track has been found in a muon chamber module and the discriminator of the backstack in front of the module has fired. The backstack requirement typically reduces the rate by a factor 3 and for the bottom chambers even by a factor 30.

The decisions from the pretrigger, muon trigger and calorimeter trigger are combined in the trigger processor to form a 32-bit word, which is checked against valid words (= trigger conditions) selected for a specific run. Of course, the pretrigger condition must be satisfied always.

Within one run (of, say 12 hours) the instantaneous luminosity can drop by more than an order of magnitude yielding different running conditions at the beginning and the end of the run. Activation of the full muon chamber area in the first level trigger at low luminosity is possible, but at high luminosity would completely saturate the system and cause a tremendous deadtime. To maintain some flexibility, several trigger areas have been defined, which can be activated or switched off quickly. The smallest one extends up to a rapidity of 1, the largest up to 2.3. The coverage in azimuthal angle of these areas has been indicated in figure 5.12 as a function of rapidity.

The first level trigger conditions for the data considered in this thesis are:

- For $L \leq 10^{30}$ cm$^{-2}$s$^{-1}$, one muon in $|η| < 1.7$ (F3) or two anywhere,
- For $L \geq 10^{30}$ cm$^{-2}$s$^{-1}$, one muon in $|η| < 0.9$ (barrel) or two anywhere,
- or one muon in F3 + a calorimeter jet with $E_T > 10$ GeV (1989 only),

where $L$ is the luminosity.
5.4.3 The Second Level Trigger

At the second level, the muon trigger employs the full drift time information. Hits in the chambers are sorted by wire number by so-called reordering memories. Wire numbers and drift times stored in these memories are read out over VME by processor boards based on the MC68020 CPU. Track finding and combining of track projections in space is performed by six of these processor boards which run in parallel. The angular resolution achieved, is approximately 4 mrad. A similar pointing requirement is used as at level 1, but with a smaller cone size (70 mrad).

Originally, a second level calorimeter trigger was also foreseen. The cancellation of the upgrade project of the electromagnetic calorimeter made its purpose useless however, so that it has not been used during these runs.

5.4.4 The Third Level Trigger

The third level trigger, which receives the complete event information over VME from the event builder, consists of 12 IBM-3081 emulators running in parallel. At this level the CD information is processed for the first time. The transportation of the large amount of data and the complicated track finding and track reconstruction programs require a processing time of 1 second per event. Crude matching algorithms are applied to reconstructed tracks in the CD and in the muon chambers to assign momenta to muon candidates. The trigger then simply selects events containing one or more high \( p_T \) muons. The \( p_T \) threshold is adjusted such, that the special event rate is about a factor 20 lower than the normal one.

5.5 Luminosity in the 1988 and 1989 Runs

The fast scintillation hodoscopes in the pretrigger setup, mentioned in section 5.4, are also used to monitor the luminosity. The latter is inversely proportional to the \( p\bar{p} \) total cross section times the acceptance of the counters: \( \sigma \cdot \varepsilon \), which can be expanded as

\[
\sigma \cdot \varepsilon = \sigma_{el} \cdot \varepsilon_{el} + \sigma_{ND} \cdot \varepsilon_{ND} + \sigma_{SD} \cdot \varepsilon_{SD} + \sigma_{DD} \cdot \varepsilon_{DD}.
\]

(5.5.1)

corresponding to the elastic and inelastic (non-diffractive, single-diffractive and double-diffractive) components of the total cross section. Elastic processes play no role in the experiment: \( \varepsilon_{el} = 0 \). The pretrigger requires a coincidence between a machine crossing strobed OR of all counter elements on the proton side of the detector and the same on the antiproton side. It is therefore sensitive to non-single-diffractive (NSD) inelastic processes for which the cross section is given by

\[
\sigma_{NSD} \equiv \sigma_{ND} + \sigma_{DD} = \sigma_{tot} - \sigma_{el} - \sigma_{SD}.
\]

(5.5.2)

Thus, the part of the \( p\bar{p} \) total cross section responsible for pretrigger rates can be derived from the measured (by UA4 and UA5) total, elastic and single diffractive \( p\bar{p} \) cross sections.
At $\sqrt{s} = 630$ GeV, they are given by\textsuperscript{91}

$$
\begin{align*}
\sigma_{\text{tot}} &= 62.8 \pm 1.5 \pm 1.5 \text{ mb}, \\
\sigma_{\text{el}} &= 13.50 \pm 0.58 \text{ mb}, \\
\sigma_{\text{SD}} &= 9 \pm 1 \text{ mb}, \\
\sigma_{\text{DD}} &= 6 \pm 1 \text{ mb}.
\end{align*}
$$

These values correspond to a non-single-diffractive cross section of: $\sigma_{\text{NSD}} = 40.3$ mb. It should be noted, that a small fraction of high-mass single-diffractive events yielding secondaries in both rapidity hemispheres can also fire the pretrigger, thus slightly increasing the effective cross section.

The pretrigger acceptance has been calculated\textsuperscript{83,84} using machine crossing data (a machine crossing trigger employs pick-up electrodes close to the beam on either side of the experiment which detect the electromagnetic field of a passing (anti)proton bunch), by feeding the sample through an offline selection of inelastic $p\bar{p}$ events and computing the fraction of passing events for which the pretrigger had fired. The acceptance is a function of the luminosity, because with increasing luminosity multiple interactions become more likely thus increasing the probability of firing the pretrigger. Theoretically, one can distinguish multiple interactions on the basis of possible combinations of the four $p\bar{p}$ processes (el, ND, SD, DD). In practice, however, it is not possible to identify all processes unambiguously. It is only possible to roughly distinguish single from double or triple interactions.
by counting the number of vertices in the CD. Therefore, the acceptance is defined as: $A = \varepsilon_1 w_1 + \varepsilon_2 w_2 + \ldots$, where $\varepsilon_n$ is the efficiency of the pretrigger to fire on an event with $n$ interactions, and $w_n$ is the luminosity dependent fraction of interaction events with exactly $n$ interactions (at $L = 1.5 \times 10^{30} \text{ cm}^{-2}\text{s}^{-1}$: $w_1 = 88\%$, $w_2 = 11\%$ and $w_3 = 1\%$).

In the presence of multiple interactions the luminosity is related to the pretrigger rate, $R$, according to

$$L = \frac{\ln(1 - \Delta t \times R)}{\Delta t \times \sigma}, \quad (5.5.4)$$

where $\Delta t = 3.8 \mu s$ is the bunch crossing time, and $\sigma$ is the effective cross section corrected for the pretrigger acceptance. Since the latter depends on $L$, the luminosity determination is an iterative procedure. Figure 5.13 shows the luminosity delivered to the UA1 experiment in the 1988 and 1989 runs. The one recorded on cassette is much lower due to the running efficiency of the detector (45% in 1988 and 75% in 1989). In total, 4.90 pb$^{-1}$ has been accumulated of which 4.66 pb$^{-1}$ corresponds to muon triggers. The systematic uncertainty on these numbers is approximately 8%, which comprises a 7% error on the effective cross section and a 3% error on the resolution of two vertices in the CD.
6 Event Reconstruction and Simulation
6. Event Reconstruction and Simulation
6.1 Introduction

In the previous chapter, we described the UA1 detector and briefly outlined how the information corresponding to complete events is recorded onto cassettes. At this stage, the so-called ‘raw’ data consists of mere sets of pulse heights, drift times, wire numbers, ADC counts, etc. To make the data suited to further analysis the cassettes are processed offline, where dedicated reconstruction programs convert the raw information into physical quantities, such as e.g. track momenta and energies in calorimeter cells. Here we consider those aspects of the event reconstruction which are relevant to the W/Z analysis presented in the next chapter. In order to allow an interpretation of the data in terms of theoretical models, it is required to perform event simulations. This will be discussed in section 6.3.

6.2 Event Reconstruction

6.2.1 Track Reconstruction in the Central Detector

Event reconstruction in the CD comprises track finding to identify charged particles and track fitting to evaluate their momenta. Track finding is performed in the xy-plane, which is both the drift plane in the CD and the bending plane of the magnet. In a first stage, strings of neighbouring hits in the individual drift volumes are chained together using parabolic paths. Subsequently, track segments of different drift volumes are combined in order to form complete tracks. The three dimensional reconstruction of these tracks is completed by including the charge division information, which yields the z-coordinate of the individual hits.

In the track fitting procedure the trajectories are assumed to have a helical form. Since the drift time related coordinates xy are measured with a much larger accuracy than the charge division coordinate z, the helix parametrization is done in projections. The radius of curvature \( R_{xy} \) is obtained from a circle fit to the track in the bending plane, while a linear fit to the track along the field yields the dip angle \( \lambda \) of the helix. Once these parameters have been established, the momentum \( p \) of the particle which created the track, is given by

\[
p = \frac{0.3 B R_{xy}}{\cos \lambda},
\]

(6.2.1)

where \( B \) is the magnetic field strength (normally 0.7 Tesla), \( R_{xy} \) is in meters and \( p \) in units of GeV.

When the parameters of all tracks have been found, the main interaction point or ‘primary vertex’ is located using the known tracks. The coordinates of the interaction point in the transverse plane (yz) are rather well known from the position of the beam (to within 200 \( \mu \)m). The coordinate along the beam line (x), however, is largely undetermined due to a bunch length of about 30 cm.

Every track is extrapolated towards the beam line in the xy-plane. Subsequently, the crossing points or points of closest approach are evaluated and entered into a histogram. Those tracks which form clusters, are studied. Usually, the largest cluster corresponds to the main vertex. A least squares fit procedure which minimizes the squared distances of the tracks to a common point, yields the vertex coordinates in the xy-plane with error estimates.
Finally, tracks which are associated to the vertex within $3\sigma$, are refitted using the vertex as an additional point. This is especially important for the linear fit along the direction of the magnetic field, because the $z$-coordinate of the main vertex is well known from the beam position, while the $z$-coordinates of a track are obtained from the charge division with an intrinsically low resolution.

6.2.2 Muon Identification

Muons are minimum ionizing particles and have a long lifetime ($c\tau = 659$ m). This combination of properties allows them to traverse large amounts of material without being stopped, a feature which forms the basis for muon identification at UA1. From all charged particles produced in the interaction point, only muons with a momentum above some threshold value ($\sim 3$ GeV) are capable of reaching the muon chambers (muons with momenta below threshold are absorbed and are said to ‘range out’). All other particles are absorbed in the calorimeter and the iron shielding behind it. Therefore, the signature for a muon consists of a track in the CD pointing to a set of hits in the muon chambers, possibly accompanied by a small amount of energy deposited in the calorimeter.

![Figure 6.1: Matching between a CD-track and a track in the muon chambers](image)

In more detail, muon identification proceeds as follows. First of all, tracks are reconstructed in the muon chambers in both projections (see figure 5.9). In the absence of delta rays, electronic double pulses, etc., each projection contains at most four hits which lie on a straight line, since the muon chambers are located outside the magnetic field. In practice, a minimum of three hits is required. Two-dimensional straight line fits determine the track parameters in each projection, which are subsequently combined to yield a single track in three dimensions. With a point resolution of approximately 0.6 mm, an angular resolution of about 3 mrad is obtained. The bottom chambers form an exception. They contain only one set of parallel drift tubes and therefore allow only one track projection to be measured. The other projection can be obtained with a much lower accuracy, however, by recording signals on both ends of the drift tubes. Differences in drift time yield the coordinate along the tube with a precision of about 0.3 m.

After tracks have been reconstructed in both the CD and the muon chambers, they are geometrically associated by a matching procedure which is illustrated in figure 6.1. All tracks
in the CD with a transverse momentum exceeding 3 GeV are extrapolated towards the muon chambers, while correctly taking into account multiple scattering of a minimum ionizing particle and bending in the magnetic field present in the calorimeter and part of the iron shielding. The average momentum loss is calculated in small steps during extrapolation. The procedure stops if the momentum of the particle drops below a threshold value of 100 MeV or if the centre of a muon chamber module is reached. In the latter case, the position and direction of the extrapolated track can be compared or 'matched' to the ones of the track reconstructed in the muon chambers. If the matching fulfills some quality requirements, to be discussed in the next chapter, a muon has been found of which the momentum is determined by the track in the CD.

6.2.3 Identification of Neutrinos and Jets in the Calorimeter

The identification of neutrinos and jets and also the measurement of the energy flow surrounding a muon employs the granularity of the UA1 calorimeter. This feature allows the definition of energy vectors as illustrated in figure 6.2: for calorimeter cell $i$ with energy $E_i$ an energy vector is defined as $\mathbf{E}_i = E_i \mathbf{u}_i$, where $\mathbf{u}_i$ is a unit vector pointing from the nominal interaction point to the reconstructed position of the calorimeter hit. If a muon is present in the event, the expected muon energy deposition is removed from the calorimeter cells before the energy vectors are reconstructed.

![Figure 6.2: Definition of an energy vector.](image)

In principle, the existence of neutrinos or non-interacting particles in general can be inferred from the energy measurement of all detected particles. In $p\bar{p}$ collisions, however, a large fraction of energy is carried off by particles with high momentum and low transverse momentum, which escape detection by remaining in the beam pipe. For this reason only the transverse component of the energy imbalance can be measured, as will be outlined below. From energy conservation in the transverse plane it follows that

$$\sum_{i=1}^{n} E_T^i = 0,$$

(6.2.2)

where the sum extends over all particles in the event. Since a neutrino is a non-interacting particle, its occurrence in an event will signal an apparent transverse energy imbalance, or missing energy: $E_T^{\text{miss}} = -E_T$. Some caution must be taken in the case of events with muons, since from a calorimetric point of view a muon also has to be largely considered as a
non-interacting particle. However, the muon detection and momentum measurement by the CD still allow the application of the transverse energy balance method even in the presence of muons. In muon events, the transverse momentum of the neutrino is defined as

\[ \mathbf{p}_T^\nu = \mathbf{E}_T^{\text{miss}} = -\mathbf{p}_T^\mu - \mathbf{E}_T, \]  

(6.2.3)

where \( \mathbf{p}_T^\mu \) is the muon transverse momentum vector measured in the CD. The muon energy deposition in the calorimeter (~2 GeV) has been subtracted from \( \mathbf{E}_T \).

In the calorimeter, a jet, which consists of a collimated shower of particles, shows up as a local cluster of energy. Several cluster finding algorithms exist\(^94\), and all make use of energy vectors such as defined above. The one applied at UA1 works as follows:

- all calorimeter cells with a transverse energy exceeding 1.5 GeV, so-called initiators, are arranged in a list of descending order in transverse energy,
- starting with the highest initiator in the list, others are associated which lie within a cone of \( \Delta R = 1.0 \) around it,
- the procedure is repeated with the next highest initiator left, until all of them have been used,
- remaining cells with an energy above 100 MeV are added to a cluster of initiators if they lie within a cone of \( \Delta R = 1.0 \) around the cluster axis.

Under the assumption that the particles which deposited the energy in the calorimeter are very light (this is reasonable, because most particles in jets are pions), one can convert the energy vectors into four-vectors (with \( p = E \)). A sum over the four-vectors of cells in a cluster then defines the four-vector of the jet.

The basic parameters of the algorithm are the threshold energy applied to the initiators and the cone size used for association. The choice of the initiator threshold is dictated by two competing requirements: a reasonable efficiency for jets of low transverse energy requires a low threshold, whereas a good rejection against fake jets from, for instance, underlying event fluctuations or instrumental problems such as hot photomultipliers, requires a high threshold. The choice of 1.5 GeV represents an optimization. The choice of the cone size is related to the ability to resolve nearby jets. Since jets in the events which are considered in this thesis, are well separated, we will not address this point further.

A similar jet finding algorithm works on tracks reconstructed in the CD. However, it should be kept in mind that in those jets the contribution from neutral particles is neglected.

### 6.2.4 Event Scanning

The full reconstruction details of an event can be inspected on a MEGATEK graphics workstation. By online manipulation (rotation, translation, zooming in and out, etc.) of the event on display, the scanner obtains a three-dimensional impression of the distribution of tracks in the CD and in the muon chambers and of energy depictions in the calorimeter. Details of measurements can be investigated and parts of the reconstruction can be redone.
It has become a standard procedure for low statistics muon samples, such as the one to be described in the next chapter, to check whether CD tracks of muon candidates show an abrupt change in curvature (a ‘kink’) by fitting parts of a track separately. This has proven to be an important tool in understanding a background to prompt muons. In fact, many offline procedures used to reject background are coded versions of strategies developed at the MEGATEK.

Figure 6.3 is a MEGATEK picture of a $Z^0 \rightarrow \mu^+\mu^-$ candidate event. It shows the cylindrical contours of the CD in the centre, the outlines of the calorimeter and the two muon chamber modules that have been hit. Only those calorimeter cells are shown, which have a reconstructed transverse energy larger than 1 GeV. The $p_T$ threshold applied to CD tracks was 0.5 GeV. Extrapolations of CD tracks of muon candidates have been indicated by dashed lines, tracks reconstructed in the muon chambers by arrows. For the lower muon, the energy deposition in the calorimeter is visible. For the other one, the transverse energy was less than 1 GeV. The CD tracks in the forward directions are beam jets originating from spectator partons in the proton and antiproton. The transverse momenta of the muons in this event are 86 GeV and 23 GeV. The mass of the $Z^0$ candidate is 98.6 GeV.
6. Event Reconstruction and Simulation

6.3 Event Simulation

6.3.1 Introduction

Because of the complex nature of p+p collisions, the response of the UA1 detector, and the reconstruction and analysis programs, one cannot directly compare the experimental results with purely theoretical calculations. To make a comparison possible, either the data have to be corrected for imperfections of the apparatus and inefficiencies in the event selection, or a detailed simulation has to be carried out of the relevant physical processes and the specific detector behaviour. For the analysis of the UA1 data, we have used the second option.

Basically, three stages can be distinguished in the simulation of an event. The first step comprises all the physics involved with a p+p collision and results in a set of particles emerging from the interaction point. This is accomplished by so-called ‘event generators’ of which several types are on the market\textsuperscript{95,96,97}. In the subsequent phase, particles are tracked through the apparatus while their interaction with detector material is simulated. Unstable particles are allowed to decay. Finally, the response to these interactions is computed for the individual detector components and a trigger simulation is performed. Events thus generated are fed through the same reconstruction programs as described above, and can be analyzed in the same way as the data. Since most of the simulation involves statistical processes (interactions between particles) the calculations are based on the Monte Carlo method\textsuperscript{98}.

6.3.2 Simulation of Weak IVB Production

The event generator most commonly used in UA1 is the ISAJET Monte Carlo program\textsuperscript{4}, which provides for the simulation of Minimum Bias\textsuperscript{99}, QCD and Drell-Yan (W$^\pm$, Z$^0$ and $\gamma^*$) events in p+p collisions. Over the years, it has been tuned to agree with observed features of the data in the experiment\textsuperscript{100}.

ISAJET uses the Drell-Yan model for IVB production. It includes the lowest order subprocess: q + $\bar{q}$ $\rightarrow$ V, and the O($\alpha_S$) subprocesses: q + $\bar{q}$ $\rightarrow$ V + g and g + q $\rightarrow$ V + q. At Born level, IVBs are produced with zero transverse momentum (p$_T$) or at most 300 MeV owing to the Fermi motion of the scattering partons in the (anti)proton. Transverse momenta are generated by the O($\alpha_S$) processes, where the radiated quark or gluon recoils against the IVB in the plane transverse to the beam. These higher order contributions can be reliably calculated for large p$_T$, but diverge for p$_T$ $\rightarrow$ 0. Therefore, a parametrization with an appropriate low p$_T$ cutoff is introduced in order to reproduce the p$_T$ behaviour predicted by Altarelli et al.\textsuperscript{101} shown in figure 6.4. The parametrization is normalized to get an integrated cross section close to the measured one.

The O($\alpha_S$) processes yield IVBs in the presence of at most one jet. Matrix elements for higher order QCD subprocesses, giving rise to multi-jet events, have not been incorporated in ISAJET. Instead, multi-jet events are obtained in the branching approximation\textsuperscript{102} as follows: the primary parton from the O($\alpha_S$) subprocess is assigned a virtual mass of order p$_T$ and then evolved using the basic QCD branching processes (q $\rightarrow$ qg, g $\rightarrow$ gg and g $\rightarrow$ q$\bar{q}$); partons generated during the evolution acquire successively lower virtual masses and the evolution stops when the virtual mass drops below a fixed cutoff value of about 6 GeV.
Below this value the hadronization is described in a non-perturbative independent fragmentation model\textsuperscript{103}. For the parton densities in the proton and antiproton, ISAJET uses non-scaling structure functions. In the present analysis, we have used the parametrization of Eichten et al.\textsuperscript{104}, solution 1.

Once the weak IVB\text{s} have been generated, their decay is described by a pure V–A current for W's, and the appropriate mix of vector and axial-vector currents for Z\textsuperscript{0}'s. The weak mixing angle used in ISAJET corresponds to $\sin^2\theta_W = 0.214$, while the vector bosons have been assigned the following masses: $M_W = 83.4$ GeV and $M_Z = 94.1$ GeV. Only three IVB decay channels contribute significantly to muon production. They are the direct decays: $W \rightarrow \mu\nu$ and $Z^0 \rightarrow \mu^+\mu^-$, and the indirect one: $W \rightarrow \tau\nu_{\tau}$ followed by $\tau \rightarrow \mu\nu_{\mu}\nu_{\tau}$. For the $\tau$ decay, ISAJET uses a branching ratio of 18\textsuperscript{\%}\textsuperscript{105}. Although this branching ratio is substantial, most muons thus produced have only a small transverse momentum owing to the three-body decay of the $\tau$. This is illustrated in figure 6.5, where the $p_T$ spectrum of muons from $\tau$'s is compared to the one of muons originating from W's directly.

---

\textbf{Figure 6.4:} Shape of the $p_T$ distribution for W and Z bosons at $\sqrt{s} = 630$ GeV\textsuperscript{101}

\textbf{Figure 6.5:} Transverse momentum of muons in $W \rightarrow \mu\nu$ (solid) and $W \rightarrow \tau\nu$ (dashed) events as predicted by ISAJET. For the decay: $\tau \rightarrow \mu\nu_{\mu}\nu_{\tau}$, a branching ratio of 18\textsuperscript{\%} was taken.
Events are completed by the generation of an 'underlying event', that is, the activity associated with spectator partons not involved in the hard scattering process. They are handled by a mechanism similar to that for producing non-diffractive Minimum Bias events at SpS collider energies\textsuperscript{106}. The activity around muons from weak IVBs is almost entirely due to the underlying event. They are said to be 'isolated' as opposed to muons produced inside jets. Since the muon isolation will be used as an IVB selection criterion, a correct simulation of the spectator parton behaviour is required. Parameters which control the spectator activity in the Monte Carlo program have been adjusted to reproduce the particle multiplicity, transverse momentum and transverse energy flow around muons and jets in muon-jet events observed in earlier UA1 data. This is illustrated in figure 6.6, where the histograms represent the data and the black dots the ISAJET prediction (including a detector simulation). Figures a), b) and c) show the charged multiplicity, transverse momentum and transverse energy flow in the same azimuthal hemisphere as the muon measured as a function of the pseudorapidity difference between the muon and the charged track (or energy.
vector). Figures d), e) and f) show similar profiles, but measured with respect to the highest \( E_T \) jet in the event pointing away from the muon (\( \Delta R(\mu, \text{jet}) > 1 \)). In fact, these plots illustrate more than a correct description of the underlying event, which is dominant in the regions: \( 1 < |\Delta \eta| < 2 \). They also show that the branching approximation and the fragmentation model used by ISAJET lead to a correct description of the observed features of jets in the region: \( |\Delta \eta| < 1 \).

A large ISAJET production of W and Z events was performed. In total, the equivalent of 50 pb\(^{-1}\) of Monte Carlo data was generated, which represents about 10 times the size of the 1988—1989 data sample. All events were subsequently fed through a simulation of the UA1 detector. A comparison of these events before and after detector simulation shows how very characteristic features of W and Z events are washed out by the resolution of the apparatus. As an example we show the transverse mass distribution of W’s before and after detector simulation (with an \( M_T \) cut at 30 GeV) in figure 6.7. The Jacobian peak collapses after detector smearing into a distribution where many W’s are reconstructed at low \( M_T \). This effect significantly complicates the analysis to be discussed in the next chapter.

![Graph](image)

Figure 6.7: ISAJET prediction for the W transverse mass before (solid) and after (dashed) detector simulation.

6.3.3 Simulation of Other Processes

Several other processes responsible for muon production in p\( \bar{p} \) collisions have been generated with ISAJET in addition to the W and Z samples mentioned above. The relevant ones are: (i) Drell-Yan production: p\( \bar{p} \rightarrow \gamma^*X \rightarrow \mu^+\mu^-X \), and (ii) heavy-flavour production: p\( \bar{p} \rightarrow c\bar{c}(b\bar{b})X \) followed by a semileptonic decay of one of the heavy quarks. The differential cross sections for these processes have been derived from UA1 measurements\(^{107}\). Also in these channels, an equivalent of 50 pb\(^{-1}\) of Monte Carlo data was generated. All events have been fed through a detector simulation and the standard reconstruction programs.
7 Data Analysis
7. Data Analysis
7.1 Introduction

In this chapter, we discuss the muon data taken in the 1988 and 1989 runs. We start with an overview of the different sources of prompt muons in $p\bar{p}$ collisions and discuss some of the technical backgrounds. An earlier study\textsuperscript{108} based on data taken in 1984 and 1985 has shown, that these different sources can fully explain the observed transverse momentum distribution of inclusively produced muons.

We then present the inclusive muon selection for the 1988 and 1989 data. It forms a sort of reference sample for many analyses carried out in UA1, including the one in this thesis. Basically, it requires events with well defined muons with some minimum transverse momentum. No kinematical cuts are made at this stage.

Subsequently, W and Z events are selected from the inclusive muon event sample by a series of cuts concerning the event topology. In principle, the characteristics of W and Z events are very distinct and should allow a clear separation from other events. However, many features are washed out by detector resolution (this was illustrated for the W transverse mass distribution in figure 6.7 in the previous chapter).

While doing the analysis, we found that understanding the $W \rightarrow \mu \nu$ event sample really meant understanding the muon background from $\pi$- and K-mesons decaying in flight. Therefore, a large fraction of this chapter will be devoted to a discussion of this particular background. Two methods have been used to select W and Z events. The first employs the standard reconstruction of muons as discussed in chapter six. A special procedure is developed to predict the size and shape of distributions in kinematical variables for the $\pi$/K decay background, while the other much smaller backgrounds are obtained with ISAJET. By normalizing signal and background distributions to the ones observed in the data, the relative contribution of each can then be determined. The second method aims at explicitly eliminating the $\pi$/K decay background by extending the muon momentum measurement to include information from the muon chambers. This method has a high rejection power against $\pi$/K decays, but unfortunately is less efficient.

We conclude this chapter with a breakdown of the overall W and Z detection efficiencies.

7.2 Inclusive Muon Production

7.2.1 Introduction

Since only a small fraction of muons produced in $p\bar{p}$ collisions originates directly from W and Z decays, we consider it important to give some general overview of the different processes responsible for muon production. In addition we will discuss some technical backgrounds to muon events and show that, once this background has been removed, the measured inclusive muon $p_T$ spectrum can be explained in terms of known physical processes.
7.2.2 Sources of Muons in Proton-Antiproton Collisions

The transverse momentum spectrum of muons observed in the UA1 detector is a steeply falling distribution. It starts at about 3 GeV, the minimum energy needed to penetrate all the material in front of the muon chambers. Muons produced via an intermediate W or Z are located at the far end of the spectrum.

Over 80% of the muons with transverse momenta below 5 GeV originate from pions and kaons, which are copiously present among the fragments of quark and gluon jets. This muon source is considered as background and will be discussed in the next subsection.

With increasing muon transverse momentum, other processes slowly gain in significance with respect to the π/K decay background. Among those, semileptonic decays of heavy flavoured mesons which arise from the fragmentation of heavy quarks (Q=c,b), form a large contribution. In principle, muons thus produced should be considered as semiprompt, because their parent mesons have mean lifetimes which correspond to several tenths of millimeters. However, since the UA1 apparatus is lacking a vertex detector, it is not possible to distinguish between main and secondary vertices at this scale.

A fraction of those heavy quarks is produced weakly via an intermediate W or Z, but most are created strongly as open QQ states by the annihilation of light quarks (q) or gluons (g): q̅q → QQ and gg → QQ. In addition, single heavy quarks are produced by the fusion of a gluon with a heavy quark from the quark sea in the (anti)proton: gQ(Q) → gQ(Q).

After quark fragmentation, the semileptonic decay of the heavy flavoured mesons yields muons via the subprocesses: b → cμν and c → sμν. These muons are surrounded by other jet fragments and hence are non-isolated. In the case of charm production, at most two muons can be produced of opposite charge, while for bottom production also like sign combinations are possible and even more muons can be produced. These mechanisms have been extensively studied by the UA1 collaboration and led to interesting results concerning mixing in the B^0 → B^0 system and the inclusive bottom cross section.

An attempt has been made to detect heavy quarks created by the weak interaction, in which two-jet events were considered where one of the jets was tagged by a muon. However, due to the overwhelming background of strongly produced heavy quarks the spectrum of the two-jet invariant mass did not show any significant enhancement near the W and Z masses.

Decays of bound states of heavy flavours such as the J/ψ and the Y, and Drell-Yan type processes can give rise to muon pairs in the final state. For such events, two classes can be naturally distinguished by their dimuon topology: (i) high mass, predominantly low p_T muon pairs, where the muons are emitted essentially back-to-back in the transverse plane (high p_T muons only for Z^0 decays), (ii) low mass, high p_T muon pairs, where the muons are emitted in a non-back-to-back configuration, sometimes with small opening angles for very low masses. Several studies have been dedicated to both classes of events.
7.2 Inclusive Muon Production

7.2.3 Backgrounds to Muon Production

In addition to muons produced in p̅p interactions, several other muon sources have to be considered which can show up as background in the analysis of muon data. With background we mean events containing muons which have not been produced in the interaction point but elsewhere in the detector or even outside the apparatus.

By far the largest background to prompt muon production comes from pions and kaons decaying in flight. This is due to a combination of facts. First of all, the branching ratio’s for the decays: \( \pi^+ \rightarrow \mu^+\nu_\mu \) and \( K^+ \rightarrow \mu^+\nu_\mu \) are 99.99\% and 63.51\% respectively. In addition, pions and kaons are practically always present among the fragments of quark and gluon jets. Since the production cross sections for light quarks and gluons in p̅p collisions are orders of magnitude higher than the ones of processes mentioned earlier, this source would make up 99\% of the muon spectrum, if all pions and kaons would decay before reaching the muon chambers. However, their lifetimes are rather long (\( c\tau(\pi^+) = 370.4 \) cm and \( c\tau(K^+) = 780.4 \) cm), so that most of these mesons have been absorbed in the calorimeter before they actually decay.

The simple geometry of the CD (a cylindrical volume) allows to make a hand calculation of the probability for mesons to decay inside the CD. Suppose \( N_0 \) mesons of a certain type with decay length \( \lambda \) are produced in a p̅p interaction. Then after having travelled a distance \( x \), there are \( N(x) = N_0\exp(-x/\lambda) \). If we define \( \Pr \) to be the probability for a decay, then the probability for no decay is given by: \( 1 - \Pr = N(x)/N_0 = \exp(-x/\lambda) \). For \( x \ll \lambda \) we can expand the exponential, so that the probability becomes: \( \Pr \approx x/\lambda \). With the decay length in the LAB frame given by: \( \lambda = (p/mc)\tau \), the probability for a meson to have decayed in a distance \( x \) becomes

\[
\Pr(\text{meson decay}) = \int_0^x \frac{mc}{c\tau p} \, ds \quad x \ll \lambda. \quad (7.2.1)
\]

The distance \( x \) can be expressed in a radial distance \( \rho \) and a polar angle \( \theta \) with respect to the beam axis by: \( x = \rho/\sin\theta \). If we use the transverse momentum \( (p_\tau = p\sin\theta) \) instead of the momentum itself, substitute particle masses, and put \( \rho \) equal to the diameter of the CD, we get the following expressions for the probabilities: \( \Pr(\pi^+ \rightarrow \mu^+\nu_\mu) = 0.02/p_\tau \) and \( \Pr(K^+ \rightarrow \mu^+\nu_\mu) = 0.15/p_\tau \). Thus, the decay probabilities are inversely proportional to the transverse momenta of the mesons.

In reality, these probabilities depend on the trigger acceptance. Exact values have been computed with a single track Monte Carlo, in which many pions and kaons were generated and allowed to decay within the trigger acceptance. The following values have been obtained

\[
\Pr(\pi^+ \rightarrow \mu^+\nu_\mu) = \frac{0.023}{p_\tau} \quad (7.2.2)
\]

\[
\Pr(K^+ \rightarrow \mu^+\nu_\mu) = \frac{0.11}{p_\tau} .
\]

So far, we only considered decays inside the CD. The absence of the electromagnetic calorimeter leaves an additional 36 cm of free space for mesons to decay, which increases the decay probabilities by a factor of about 1/3.
Many decays which occur in the CD can be recognized, since the track residuals show a typical kink-like behaviour as illustrated in figure 7.1 (we will frequently use the word ‘kink’ to denominate tracks associated with π/K decays). In general pion decays are more difficult to reconstruct, because their small mass leads to muon momentum of only 30 MeV in the pion CM frame, whereas for the heavier kaons the muons get 236 MeV. If a decay remains undetected, the misreconstructed track often seems rather stiff resulting in a large overestimation of its momentum.

There are several ways to protect against this type of background. First of all, the track fit usually is of a bad quality, which is reflected in a bad $\chi^2$. In addition, extrapolation to the muon chambers is done with starting values for the track momentum and direction, which are to some extend off from the correct values. This leads to a bad matching. Therefore, the requirement of good quality and well matched tracks removes some fraction of this background. A further reduction is achieved by the application of a kink-finder, which consists of a software program that checks potential muon tracks for decay points by stepping through the track and fitting the two parts before and after the breakpoint separately\textsuperscript{112}. If the two momenta differ significantly from the original one, the track is identified as having a kink. If the track segment nearest to the interaction point has a higher momentum than the other one, the kink finder labels the muon as a decay and rejects it. This procedure can be pursued in more detail at the scan table, where more consistency checks can be made at the same time.

Figure 7.1: Residuals of the reconstructed tracks associated with a meson decay displaying a typical kink structure. The event vertex is located on the left of the picture.

A substantial flux of muons from cosmic rays hits the UA1 apparatus, which is shielded by only five meters of earth. Although this is useful for calibration purposes, it
might form a potential background to muon events. Fortunately, the muon flux which passes through the interaction region is much smaller, even more so if one requires it to be exactly in time with a bunch crossing. Those cosmics which get through, are traced with a cosmic ray finder. It rejects all events in which the muon is not attached to a vertex. In case the muon is attached, it searches systematically for another track opposite to the muon in the xy-plane. If such a track is found, the vertex is removed from both the muon and the opposite track and the two are combined in a single track fit. The resulting track has to go through the vertex (within 3σ), otherwise the muon is identified as a cosmic and the event is rejected. Remaining cosmics are rejected by scanning.

In principle, hadrons can pass through the calorimeter without interacting, thus giving hits in the muon chambers. However, the probability to traverse more than 9 interaction lengths is very small. Test beam measurements have shown that the probability to penetrate the whole detector is less than 10^{-4} per incident hadron^{113}. Another possibility is, that some hadronic showers are not fully contained by the calorimeter. Escaping charged secondaries might then again give hits in the muon chambers. The same test beam measurements have indicated that the probability for this to happen is also less than 10^{-4} per incident hadron.

Since the muon identification procedure consists of the geometric matching between a reconstructed track in the CD and one in the muon chambers, it is possible that the wrong CD track has been chosen. This is indeed a problem for muons inside jets. For muons from W and Z decays, however, the association is unambiguous, because they are isolated and have a high transverse momentum.

### 7.2.4 Inclusive Muon $p_T$ Spectrum

In order to situate muons originating from W and Z decays in the inclusive muon $p_T$ spectrum, we consider an earlier UA1 measurement^{108} concerning heavy flavour production based on an inclusive sample of events containing a muon with $p_T > 6$ GeV and $|η| < 1.5$. The data sample used for this study was collected in 1984 and 1985 at a centre-of-mass energy of $\sqrt{s} = 630$ GeV and corresponds to an integrated luminosity of 556 nb$^{-1}$. The only cuts applied to the sample were quality cuts, such as the requirements of well reconstructed and well matched muon tracks, to ensure well defined muons. In addition, a kink finder was applied, to remove as much decay background as possible. Finally, a special technique was used, which will be touched upon later in this chapter, to calculate the amount of remaining decay background. This was subsequently subtracted from the inclusive sample.

A large (ISAJET) Monte Carlo production was performed for processes which contribute to the inclusive muon sample. All events were passed through the full UA1 detector simulation and reconstruction programs. They were then selected and analysed with the same programs as the data, thus allowing for a comparison of the data with theoretical predictions.

In figure 7.2 the inclusive muon cross section $dσ/dp_T$ after subtraction of the decay background is compared to the equivalent prediction for the muon cross section obtained by summing all the physics processes considered in ISAJET, namely charm and bottom production, W and Z decay, Drell-Yan, J/ψ and Υ production. The dashed-dotted curve shows the contribution from W and Z decays alone. Notice the flattening of the spectrum for
transverse momenta larger than 15 GeV, which is due to the onset of muon production via intermediate W and Z bosons.

The conclusions drawn in the same reference, which are relevant for the present analysis are: (i) the $\pi/K$ decay background fraction decreases from 75% to 20% as $p_T^{\mu}$ increases from 6 GeV to 20 GeV, (ii) after background subtraction, the contribution from W and Z decays increases with $p_T^{\mu}$ from ~10% at 15 GeV and exceeds the $c\bar{c}/b\bar{b}$ contribution for $p_T^{\mu} > 25$ GeV.

Since the runs in 1984 and 1985, the situation has slightly changed (chapter 5). Basically, the CD momentum resolution has degraded and the removal of the electromagnetic calorimeter has increased the free space available for pions and kaons to decay. Therefore, we expect an increase of decay background and probably a shift of the point of W and Z domination over $c\bar{c}/b\bar{b}$ production towards higher $p_T^{\mu}$.

![UA1](image)

Figure 7.2 The inclusive muon $p_T$ spectrum\(^{108}\) (corrected for decay background and acceptance). The solid curve shows the sum of the ISAJET predictions. The contributions from $W \rightarrow \mu\nu$ and $Z \rightarrow \mu\mu$ are shown as a separate curve (dot-dashed).
7.3 Inclusive Muon Selection

7.3.1 Introduction

The purpose of the inclusive muon selection is to provide a sample of events which contain at least one good quality muon with some minimum \( p_T \), but to which no further kinematical cuts are applied. The \( p_T \) cut is chosen such that it suits many different analyses while maintaining a manageable sample size.

Some sort of preselection of muon events was already performed online. During data taking, two data streams were maintained (see section 5.4). All events accepted by the second level trigger were passed on to a third level trigger running in the emulators. Its task was not to reject events, but to select an enriched sample of ‘special’ events. The algorithm applied, searched for a CD track with a \( p_T \) exceeding some threshold pointing to a track in the muon chambers. The threshold was chosen such, that the special data stream corresponded to about 5% of the normal data stream. This was achieved by a \( p_T \) threshold of 5 GeV. In total 24 million events were collected in 1988 and 1989. The analysis in this thesis is based on the subsample of ‘special’ events, of which there are in total: 1181348.

Clearly, the pointing requirement used by the third level trigger is a very crude one, since time consuming extrapolation software can not be applied at this stage. Moreover, the calibration constants used online are only approximate. Better ones can be used offline after detailed studies of the detector behaviour have been carried out. The inclusive selection presented here makes use of the final sets of calibration constants.

7.3.2 Definition of Some Quality Variables for the Muon Track

In order to select well measured tracks, we need some quantitative measure of the quality of a track. Besides some obvious quantities, such as track length and number of hits per track, two other variables will be introduced. One concerns the quality of the track fit and the other the quality of the matching between the CD track and the track in the muon chambers.

In section 6.2.1 it was mentioned that the drift time related coordinates \( x_y \) and the charge division coordinate \( z \) of a CD track are fitted independently. For both fits, a sum of squares is formed

\[
X^2_{(N)} = \sum_{i=1}^{N} \frac{r_i^2}{\sigma_i^2},
\]

(7.3.1)

where \( r_i \) is the point deviation from the fit, \( \sigma_i \) the corresponding point accuracy and \( N \) the number of CD hits that participate in the fit. Suppose that the CD measurements are free of systematics. In that case, the variables \( r_i/\sigma_i \) are standard normally distributed and the function \( X^2_{(N)} \) has a chi-square distribution, \( \chi^2 \), with \( N \) degrees of freedom. In the limit \( N \to \infty \) the latter approaches the normal distribution with mean \( N \) and variance \( 2N \).
The following two quantities \(^{(114)}\)

\[
g(X_{(N)}) = \frac{X^2_{(N)} - N}{\sqrt{2N}} \tag{7.3.2}
\]

and

\[
f(X_{(N)}) = \sqrt{2X^2_{(N)}} - \sqrt{2N - 1} \tag{7.3.3}
\]

are standard normal for large \(N\). The second function, \(f\), is used to define the quality of the track fits, because it approaches standard normal behaviour faster with increasing \(N\)\(^{(115)}\). In the following, we use the notation \(f_{xy}\) and \(f_z\) for the quality of the track fit in the drift time and charge division coordinate respectively.

The standard normal behaviour of the function \(f\) is spoiled, because of the systematic uncertainties in the CD measurements. This is especially true for the charge division coordinate. In order to account for slight variations of the CD performance, \(f_{xy}\) and \(f_z\) are averaged per day and corrected for a possible offset from zero.

Other measures of track quality are the number of points used in the fit and the projected track length \(l_{xy}\). Usually we require a minimum number of points which yields some rejection power against \(\pi/K\) decays, and a minimum \(l_{xy}\). A track with low \(l_{xy}\) would be nearly parallel to the magnetic field direction, and would therefore undergo little magnetic deflection; moreover, it would depend largely on the poorly measured \(z\) coordinate.

In section 6.2.2 we discussed how muons are identified by comparing the parameters of a track which has been extrapolated from the CD to the muon chambers with those of the track in the muon chambers themselves. The quality of the matching can be expressed by the following \(\chi^2\) functions

\[
\chi^2_{\text{position}} = \frac{\Delta x_1^2}{\sigma_{x_1}^2 + F(x_1)} + \frac{\Delta x_2^2}{\sigma_{x_2}^2 + F(x_2)}, \tag{7.3.4}
\]

\[
\chi^2_{\text{angle}} = \frac{\Delta \phi^2}{\sigma_{\phi}^2 + F(\phi)} + \frac{\Delta \lambda^2}{\sigma_{\lambda}^2 + F(\lambda)}, \tag{7.3.5}
\]

where \(\Delta x_1, \Delta x_2, \Delta \phi\) and \(\Delta \lambda\) represent the differences in the coordinates and angles of the extrapolated track and the muon chamber track. The associated errors \(\sigma_{x_1}, \sigma_{x_2}, \sigma_{\phi}\) and \(\sigma_{\lambda}\) comprise both the extrapolation error and the error on the muon track. To account for systematic uncertainties on the chamber alignment, these \(\chi^2\) functions are calculated with additional constant error components or ‘error floors’ represented by the \(F\) functions. In general an average of the two \(\chi^2\) functions is used.

This average \(\chi^2_{\text{avg}}\) is used in an early stage of the event selection to require a loose matching, mainly to reject some decay background. A different algorithm is applied later on, which exploits a better knowledge of the alignment between the CD and the muon chambers. This is achieved by a calibration of the matching procedure with cosmic rays. When the muon track and extrapolated CD track are compared, each of the four variables is considered individually. This provides realistic error distributions and mean effective positional offsets for each muon chamber with respect to the CD. The algorithm corrects for the shifts and expresses the significance of the matching in standard deviations \(n(\sigma_m)\) in each of its
variables. The shift corrections are applied to the apparent position of the muon chambers in the software description of the detector, even though one knows that these offsets are caused by systematic uncertainties in the CD reconstruction. However, it is the simplest solution to the problem.

7.3.3 Details of the Selection

The event selection was done in two stages. In the first, cuts were rather loose, so that knowledge of calibration constants was not critical. It was applied shortly after data had been taken for an initial reduction of the sample size. In the second stage, tighter cuts were applied which rely much more on good knowledge of the calibration constants. In fact, it was applied several times as calibrations improved with time. Besides track quality requirements, also dedicated programs were applied to reject particular types of background (such as cosmic rays, π/K decays, etc.). The selection criteria are summarized in table 7.1. The loose selection reduced the inclusive muon sample by about a factor 10, from 1181348 events to 115943 events. After the tight selection, only 41936 events remained. For these events some relevant distributions are shown in figure 7.3.

<table>
<thead>
<tr>
<th>Loose selection:</th>
<th>Tight selection:</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>IN:</strong> 1,181,348 events</td>
<td><strong>IN:</strong> 115,943 events</td>
</tr>
<tr>
<td>- The μ-track must be associated with a vertex (not necessarily the primary vertex).</td>
<td>- The μ-track must be associated with a primary vertex.</td>
</tr>
<tr>
<td>- μ-track quality: (a) # points ≥ 10, (b) f_{xy} &lt; 6.0, (c) f_{z} &lt; 9.0.</td>
<td>- μ-track quality: (a) # points ≥ 20, (b) l_{xy} &gt; 0.4 m, (c) f_{xy} &lt; 3.0, (d) f_{z} &lt; 9.0.</td>
</tr>
<tr>
<td>- The μ-track must pass the cosmic ray rejection.</td>
<td>- The μ-track must pass the cosmic ray rejection.</td>
</tr>
<tr>
<td>- p_{T}^{\mu} &gt; 8 GeV.</td>
<td>- The μ-track must pass the kink rejection.</td>
</tr>
<tr>
<td><strong>OUT:</strong> 115,943 events</td>
<td>- The μ-track must pass the leakage rejection (fiducial cuts).</td>
</tr>
<tr>
<td></td>
<td>- Tight matching: n(σ_{m}) &lt; 4 in each of the four variables.</td>
</tr>
<tr>
<td></td>
<td>- p_{T}^{\mu} &gt; 8 GeV.</td>
</tr>
<tr>
<td></td>
<td><strong>OUT:</strong> 41,936 events</td>
</tr>
</tbody>
</table>
Plot a) in figure 7.3 shows the steeply falling behaviour of the muon $p_T$ distribution between 8 GeV and 100 GeV. In all other plots two distributions are shown, one corresponding to the entire tight selection sample, the other to the subsample of events with $p_T^\mu > 12$ GeV. The pseudorapidity distribution shows a dip at $\eta = 0$. This is due to the limited acceptance of the CD in the central region for tracks along the direction of the magnetic field. The bumps for $|\eta| > 1$ reflect the step-like behaviour of the geometrical acceptance of the muon chambers which goes from $\sim 65\%$ for $|\eta| < 1$ to $\sim 85\%$ for $1 < |\eta| < 2$, and an increase of the $\pi/K$ decay background towards the forward region, where beam jets form an additional supply of pions and kaons. Both the number of CD hits and the projected track length can be used to further improve the quality of the muons. However, the first quantity allows to make a more uniform selection in azimuthal angle, while a cut on the second is
biased towards vertical tracks. Both chi-square functions in \( xy \) and in \( z \) are slightly offset from zero and have standard deviations which are a little larger than one. This is due to some systematics in the measurements, but also to the fact that the minimum number of CD hits in this selection is 20. The approximation of these functions defined in equation (7.3.3) by a standard normal distribution becomes good only for more than about 30 hits\(^9\).

### 7.3.4 Vertex Refit

Our interest in the tight selection sample concerns the subsample of high \( p_T \) muons. Since the CD momentum resolution degrades with increasing momentum (section 5.3.2: \( \Delta p/p \propto p \)), muon momenta in our sample are often badly measured. An effort was made to remove the worst of these before applying the actual \( W \) and \( Z \) selections.

To understand the method we have to anticipate the discussion on \( W \) and \( Z \) event topologies in the next section. Basically, these events have low track multiplicities and contain one (\( W^\prime \)'s) or two (\( Z^\prime s \)) isolated high \( p_T \) tracks. In section 6.2.1 it was mentioned that after the coordinates of the event vertex have been determined, each track which is associated to it is refitted with the vertex included. Generally this improves the momentum measurement, because the position of the vertex yields some additional information. However, for the typical \( W \) and \( Z \) event topologies this is not true. To see this, consider the determination of the vertex position.

We recall that it is mainly the \( x \)-coordinate of the vertex along the beam that needs to be determined. Clearly, forward tracks introduce a larger uncertainty in the \( x \)-coordinate than central ones (see figure 7.4), because a small mismeasurement in \( \theta \) leads to a large shift in \( x \). Therefore, tracks which participate in the vertex fit are weighted by a factor \( \sin\theta \). For events with low track multiplicities and one stiff central track, the \( x \)-coordinate of the vertex is largely correlated to the one of the stiff track. Therefore, the position of the vertex yields little additional information when the track is refitted with the vertex included.

The problem can be partially cured by refitting the vertex, while excluding the high \( p_T \) track from the fit. In this case, the vertex does provide additional information when the track is refitted including the new vertex. The effect on badly measured high \( p_T \) tracks is twofold. First of all, in some cases the position of the vertex has shifted so much, that the track under consideration no longer becomes associated to it within \( 3\sigma \). Since the inclusive muon selection requires the muon track to be attached to the main event vertex, this type of event will be rejected. Secondly, adding an independent point to an artificially high \( p_T \) track (e.g. unrecognized kinks of \( \pi/K \) decays) tends to increase the track curvature and hence to lower its momentum. Therefore, some of these events will be rejected in a later high \( p_T \) muon selection. The only drawback of the method is that some \( W \) and \( Z \) events have a very low
track multiplicity. For such events, the vertex position can not be determined very well, which leads to a loss of good events.

In more detail the vertex refit procedure was performed as follows. From the tight selection sample all events containing a muon with $p_T > 12$ GeV were selected (in total: 6806 events). To save time, of those events only the ones with an isolated muon ($I_p(\Delta R = 0.4) < 2$ GeV, where the isolation variable I is defined in section 7.4.2) were subject to a vertex refit. Instead of only excluding the high $p_T$ muon from the refit, all tracks with $p > 10$ GeV were excluded. We have chosen to use $p$ as a criterion rather than $p_T$, because this also removes the low weight, high momentum forward tracks from the refit. Finally, the refitted events were rejected if: (i) the muon was no longer associated to the main vertex, (ii) it had $p_T < 10$ GeV, (iii) it was considered a kink by the kink finder, or (iv) it failed the matching criteria. In total 3924 events passed this selection. They form the input sample to the W and Z selections to be discussed below. Figure 7.5 shows the associated $p_T$ spectrum.

![Transverse momentum spectrum for muons passing the vertex refit procedure.](image)

**7.4 Selection of $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ Candidates, Method 1**

**7.4.1 Introduction**

In this section, we discuss the details of the $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ event selections. An initial comparison will be made with ISAJET predictions. The next section deals with the calculation of the $\pi/K$ decay background, while all information is combined in section 7.6.

**7.4.2 $W \rightarrow \mu\nu$ Selection**

The signature for a $W \rightarrow \mu\nu$ event is given by an isolated, high $p_T$ muon accompanied by large missing transverse energy induced by a non-interacting neutrino. Usually, the muon and neutrino are produced back-to-back in the transverse plane, because of the large W mass and the fact that most W bosons are produced at low $p_T$. For the small fraction
of high $p_T$ W bosons, the muon-neutrino system is recoiling against one or more jets. If these recoil jets are absent, the only activity in the event apart from the muon comes from the spectator partons in the proton and antiproton which did not participate in the hard interaction.

The $W\rightarrow\mu\nu$ event selection is summarized in table 7.2. An initial separation from the bulk of the low $p_T$ phenomena is obtained by requiring a minimum $p_T$ of 15 GeV (see figure 7.2). By tightening the $\mu$-track quality requirements with respect to the ones in the inclusive muon selection, more badly measured tracks and kinks from $\pi/K$ decays can be removed. The cut on the projected track length remains the same, but the required number of CD hits used in the $\mu$-track fit is increased while the maximum allowed values for the chi-square functions in the residuals of the drifttime and charge division fits are decreased.

Contrary to W decays, $\pi/K$ decays and heavy flavour decays yield muons which are not isolated. The isolation can be quantified by measuring the transverse momentum or transverse energy flow inside a cone of radius $\Delta R$ around the muon. This cone is illustrated in figure 7.6. The variables $I_p$ and $I_E$ in table 7.2 are defined as follows

$$I_p(\Delta R = 0.4) = \sum_{\Delta R = 0.4} p_T,$$

$$I_E(\Delta R = 0.7) = \sum_{\Delta R = 0.7} E_T,$$

(7.4.1) (7.4.2)

where the sum in (7.4.1) runs over all CD tracks in the cone and the sum in (7.4.2) over all calorimeter cells in the cone. A larger cone size is used for the energy isolation to allow for the coarse granularity of the hadron calorimeter. The third isolation requirement in table 7.2 rejects events if they contain a jet with transverse energy exceeding 10 GeV in a cone of

<table>
<thead>
<tr>
<th>W → $\mu\nu$ Selection</th>
</tr>
</thead>
<tbody>
<tr>
<td>• $p_T^{\mu} &gt; 15$ GeV</td>
</tr>
<tr>
<td>• $\mu$-track quality:</td>
</tr>
<tr>
<td>(a) # points $\geq 30$</td>
</tr>
<tr>
<td>(b) $l_{xy} &gt; 0.4$</td>
</tr>
<tr>
<td>(c) $f_{xy} &lt; 1.0$</td>
</tr>
<tr>
<td>(d) $f_\gamma &lt; 5.0$</td>
</tr>
<tr>
<td>• $\mu$-Isolation:</td>
</tr>
<tr>
<td>(a) $I_p(\Delta R = 0.4) &lt; 1$ GeV</td>
</tr>
<tr>
<td>(b) $I_E(\Delta R = 0.7) &lt; 5$ GeV</td>
</tr>
<tr>
<td>(c) $I_f(\Delta R = 0.7) &lt; 10$ GeV</td>
</tr>
<tr>
<td>• No CD jet with $P_T &gt; 5$ GeV b-t-b to $\mu$</td>
</tr>
<tr>
<td>• $M_T(\mu, \nu) &gt; 30$ GeV</td>
</tr>
</tbody>
</table>
radius 0.7 around the muon. It aims at removing a small fraction of heavy flavour decays where the muon has a large $p_T$ relative to the jet axis and ends up outside the jet core. Since the energy flow is concentrated in the centre of the jet, such muons might appear isolated according to the first two criteria.

The next cut rejects events if they contain a jet with transverse momentum larger than 5 GeV opposite to the muon in the transverse plane, that is, if $\Delta\phi(\mu, \text{jet}) > 150^\circ$. This again removes some events with muons from heavy flavour and $\pi/K$ decays, because these are typically two-jet events (where one jet contains the muon) and momentum conservation in the transverse plane requires these jets to be back-to-back in $\phi$.

Also the last requirement in table 7.2 is an effort to cut away events with muons from heavy flavour and $\pi/K$ decays. The transverse mass is expected to be much smaller for these backgrounds, where muon and neutrino are emitted in the same direction, than for $W$ events, where muon and neutrino are emitted back-to-back in $\phi$.

The $W \rightarrow \mu\nu$ selection criteria were satisfied by 526 events. All of these were scanned on the Megatek and checked for kinks in the muon track missed by the kinkfinder. In 84 events the $\mu$-track definitely contained a kink, while for 50 events it was unclear whether there was a kink or not (table 7.3). These remaining events with definite and doubtful kinks indicate that the cuts applied in the $W \rightarrow \mu\nu$ selection were not sufficient to remove all of the $\pi/K$ decay background.

Table 7.3: Scan summary of events that pass the $W \rightarrow \mu\nu$ selection.

<table>
<thead>
<tr>
<th>$\mu$-track quality</th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>good</td>
<td>392</td>
</tr>
<tr>
<td>definite kink</td>
<td>84</td>
</tr>
<tr>
<td>doubtful kink</td>
<td>50</td>
</tr>
<tr>
<td>total</td>
<td>526</td>
</tr>
</tbody>
</table>
7.4 Selection of $W \rightarrow \mu \nu$ and $Z \rightarrow \mu^+\mu^-$ Candidates, Method 1

7.4.3 Backgrounds to $W \rightarrow \mu \nu$

The large amount of decay background is mainly due to the bad momentum resolution of the CD. A mismeasurement of the muon momentum almost automatically induces missing energy back-to-back with the muon in $\phi$ via relation 6.2.3, which can create an artificially large transverse mass. The transverse momenta of kinks are often overestimated and can thus be pushed over the threshold of 15 GeV. Therefore, the only real tool we have against this background is the muon isolation. Since the $\pi/K$ production rate is many orders of magnitude higher than the one for $W$'s, there will always be a small tail that passes the isolation criteria.

In addition to heavy flavour and $\pi/K$ decays, we expect two other sources of background. The first comes from the decay $W \rightarrow \tau v$ followed by $\tau \rightarrow \mu \nu \nu$. The three-body decay of the tau usually gives softer muons than the ones originating directly from a $W$ decay (see figure 6.5 in the previous chapter), so that most are removed by the $p_t^\tau$ cut. The other background comes from the Drell-Yan processes $\gamma^* \rightarrow \mu^+\mu^-$ and $Z \rightarrow \mu^+\mu^-$, with one of the muons missed in the reconstruction. This can happen in the horizontal plane in the central region of the CD, where the acceptance for track reconstruction is worst. Such events look very similar to $W \rightarrow \mu \nu$ events, because the reconstructed muon is isolated, the unseen muon plays the role of the neutrino and the two originate from a high mass particle. The contributions from these backgrounds were calculated with ISAJET. The results are listed in table 7.4. For the background from tau's, universality was assumed ($\text{Br}(W \rightarrow e\nu) = \text{Br}(W \rightarrow \mu\nu) = \text{Br}(W \rightarrow \tau\nu)$), the branching ratio for the decay $\tau \rightarrow \mu \nu \nu$ was taken to be 18%\textsuperscript{116} and the $W$ production cross section times branching ratio into $\tau\nu$ was normalized to a result from UA2\textsuperscript{117} measured in the electron channel: $660 \pm 15 \pm 37$ pb. Using the same normalization for $W \rightarrow \mu \nu$ yields $344 \pm 21$ events, so that the measured fraction of muons from $W \rightarrow \tau\nu$ with respect to all muons coming from $W$ decays is 5% after all cuts.

If we add signal and backgrounds under the assumptions just mentioned, we predict in total: $394 \pm 22$ events. This compares well with the observed number of events containing good muon tracks given in table 7.3. However, it is clearly not desirable to reject events with kinks on a scanning basis due to the high level of subjectivity of the scanner. Therefore, they either have to be explained by an additional background calculation or they have to be removed by software. Both ways will be pursued later in this chapter.

<table>
<thead>
<tr>
<th>process</th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>$W \rightarrow \tau\nu$, $\tau \rightarrow \mu\nu\nu$\textsuperscript{†}</td>
<td>$18 \pm 2$</td>
</tr>
<tr>
<td>$Z \rightarrow \mu^+\mu^-$</td>
<td>$9 \pm 1$</td>
</tr>
<tr>
<td>$\gamma^* \rightarrow \mu^+\mu^-$</td>
<td>$16 \pm 6$</td>
</tr>
<tr>
<td>$c\bar{c}/b\bar{b}$ decays</td>
<td>$7 \pm 4$</td>
</tr>
<tr>
<td>Total</td>
<td>$50 \pm 8$</td>
</tr>
</tbody>
</table>

\textsuperscript{†} Assuming $\sigma_W \text{Br}(W \rightarrow \tau\nu) = 660\pm15\pm37$ pb, universality and $\text{Br}(\tau \rightarrow \mu\nu\nu) = 18\%$
7.4.4 \( Z \to \mu^+\mu^- \) Selection

The process \( Z \to \mu^+\mu^- \) looks in many respects very similar to the process \( W \to \mu\nu \). The second muon, however, can be directly reconstructed whereas the neutrino from a \( W \) decay has to be inferred from other measurements. It proves to be an important tool to discriminate against backgrounds.

The cuts on the first muon in the \( Z \) selection (table 7.5) are practically the same as the ones in the \( W \) selection. The only differences are the cut on the number of CD hits, which is slightly relaxed, and the absence of a veto of events with a jet back-to-back to the muon. The second muon is identified as a second CD track with \( p_T > 10 \) GeV which must be loosely isolated; it may have poorer track quality. If the CD track extrapolates to an active area of the muon chambers, an external muon track is required in addition. Both muons must be associated with the same vertex and, finally, a cut is made requiring a dimuon mass of \( M(\mu_1,\mu_2) > 50 \) GeV. In total 60 events fulfilled the selection requirements, which were validated on the Megatek. None of the muon tracks had a kink.

Table 7.5: Details of the \( Z \to \mu^+\mu^- \) selection criteria.

<table>
<thead>
<tr>
<th>( Z \to \mu^+\mu^- ) Selection</th>
<th>First muon</th>
<th>Second muon</th>
</tr>
</thead>
<tbody>
<tr>
<td>• ( p_T^\mu &gt; 15 ) GeV</td>
<td>• ( p_T^\mu &gt; 10 ) GeV</td>
<td></td>
</tr>
<tr>
<td>• ( \mu )-track quality: (a) # points ( \geq 25 )</td>
<td>• ( \mu )-track quality: (a) # points ( \geq 15 )</td>
<td></td>
</tr>
<tr>
<td>(b) ( l_{xy} &gt; 0.4 )</td>
<td>(b) ( f_{xy} &lt; 5.0 )</td>
<td></td>
</tr>
<tr>
<td>(c) ( f_{xy} &lt; 1.0 )</td>
<td>(c) ( f_z &lt; 5.0 )</td>
<td></td>
</tr>
<tr>
<td>(d) ( f_z &lt; 5.0 )</td>
<td></td>
<td></td>
</tr>
<tr>
<td>• ( \mu )-Isolation: (a) ( I_P(\Delta R=0.4) &lt; 1 ) GeV</td>
<td>• ( \mu )-Isolation: (a) ( I_P(\Delta R=0.4) &lt; 10 ) GeV</td>
<td></td>
</tr>
<tr>
<td>(b) ( I_E(\Delta R=0.7) &lt; 5 ) GeV</td>
<td>(b) ( I_E(\Delta R=0.7) &lt; 10 ) GeV</td>
<td></td>
</tr>
<tr>
<td>(c) ( I_I(\Delta R=1.0) &lt; 10 ) GeV</td>
<td></td>
<td></td>
</tr>
<tr>
<td>• Both muons must be attached to the same vertex</td>
<td></td>
<td></td>
</tr>
<tr>
<td>• ( M(\mu_1,\mu_2) &gt; 50 ) GeV</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

The same ISAJET events as mentioned above were passed through the \( Z \to \mu^+\mu^- \) selection to compute the size of the backgrounds. The only non-negligible background appears to come from Drell-Yan production of two muons via a virtual photon (table 7.6). The decay \( Z \to \tau^+\tau^- \) followed by \( \tau \to \mu\nu\nu \) does not contribute, because the 18% branching ratio comes in twice and the muons are relatively soft. The requirement of a second track with \( p_T > 10 \) GeV effectively eliminates all background from heavy flavour decays. A calculation in the next section shows that also the contribution from \( \pi/K \) decays is negligible.
Table 7.6: $Z \rightarrow \mu^+\mu^-$ candidates and background.

<table>
<thead>
<tr>
<th></th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>Candidates</td>
<td>60</td>
</tr>
<tr>
<td>$\gamma^* \rightarrow \mu^+\mu^-$ background</td>
<td>2 ± 1</td>
</tr>
<tr>
<td>$Z \rightarrow \mu^+\mu^-$ observed</td>
<td>58 ± 7.7 ± 1</td>
</tr>
</tbody>
</table>

Figure 7.7 shows the mass distribution of unlike sign dimuons with a mass cut at 30 GeV for the data after subtraction of the Drell-Yan background and the ISAJET prediction for $Z \rightarrow \mu^+\mu^-$ after detector simulation and reconstruction. The observed width of the $Z^0$ peak is well reproduced by our Monte Carlo calculation and is dominated by the momentum measurement error. Also the observed muon transverse momentum and muon isolation for both $Z^0$ legs shown in figure 7.8 are well reproduced by our Monte Carlo simulation.

Figure 7.7: Mass of unlike sign dimuons with invariant mass larger than 30 GeV for the data after background subtraction (data points) and Monte Carlo (histogram).
7.5 Evaluation of the $\pi$/K Decay Background

7.5.1 The Method

Earlier in this chapter we discussed the background from pion and kaon decays in flight to prompt muon production. In this section we explicitly compute its contribution to the W and Z samples described in the previous section. The technique differs from the other background calculations, because real data is used instead of events generated by Monte Carlo. We first give an overview of the method and then present more details of the calculation.

The aim of the method is to create an event sample representative of real decay background events, from which distributions in several kinematical variables can be extracted. Passing this sample through the W/Z selection then yields an estimate of the size of the background.

First, consider the $p_T$ spectrum of muons from $\pi$/K decays, $d\sigma^{\mu}/dp_T^\mu$. If one assumes the transverse momentum of the muon to be the same as the one of its parent pion or kaon, that is, $p_T^\mu = p_T^{\pi(K)} \equiv p_T$, then the muon spectrum can be expressed in its parent spectra as follows

$$\frac{d\sigma^{\mu}}{dp_T} = \frac{d\sigma^{\pi}}{dp_T} P_\pi(p_T) + \frac{d\sigma^K}{dp_T} P_K(p_T), \quad (7.5.1)$$

where $d\sigma^{\pi(K)}/dp_T$ is the $\pi(K)$ $p_T$ spectrum and $P_\pi(K)(p_T)$ the $\pi(K)$ decay probability. However, the pion and kaon transverse momenta are not equal to the ones of their daughter muons, but merely represent upper limits. Therefore, relation (7.5.1) should be modified
according to

\[
\frac{d\sigma^\mu}{dp_T^\mu} = \int dp_T^\pi F_\pi(p_T^\pi, p_T^\mu) \frac{d\sigma^\pi}{dp_T^\pi} + \int dp_T^K F_K(p_T^K, p_T^\mu) \frac{d\sigma^K}{dp_T^K},
\]

(7.5.2)

where the functions \( F_{\pi(K)}(p_T^{\pi(K)}, p_T^{\mu}) \) are probability densities which yield for each possible \( p_T^{\pi(K)} \) the \( \pi(K) \) decay probability as a function of \( p_T^{\mu} \). The \( \pi(K) \) \( p_T \) distribution can be obtained from the inclusive charged hadron \( p_T \) spectrum, \( d\sigma^h/dp_T^h \), and measurements by UA2\textsuperscript{118} and UA5\textsuperscript{119} of the relative \( \pi(K) \) fraction, \( f_{\pi(K)}(p_T^{\pi(K)}) \). This gives

\[
\frac{d\sigma^\mu}{dp_T^\mu} = \int dp_T^h \left\{ F_\pi(p_T^h, p_T^{\mu}) f_\pi(p_T^h) \frac{d\sigma^h}{dp_T^h} + F_K(p_T^h, p_T^{\mu}) f_K(p_T^h) \frac{d\sigma^h}{dp_T^h} \right\},
\]

(7.5.3)

which can be written in a more compact form, if we define a new probability density for hadrons as

\[
F(p_T^h, p_T^{\mu}) = F_\pi(p_T^h, p_T^{\mu}) f_\pi(p_T^h) + F_K(p_T^h, p_T^{\mu}) f_K(p_T^h).
\]

(7.5.4)

Equation (7.5.3) then becomes

\[
\frac{d\sigma^\mu}{dp_T^\mu} = \int dp_T^h F(p_T^h, p_T^{\mu}) \frac{d\sigma^h}{dp_T^h},
\]

(7.5.5)

where the probability density has to be computed (subsection 7.5.2) and the hadron \( p_T \) spectrum can be measured (subsection 7.5.3).

The method can be extended to yield distributions in other kinematical variables by treating the hadron \( p_T \) spectrum on an event-by-event basis, so that the topological details of each event are available. In this case, one replaces the hadron with transverse momentum \( p_T^h \) by a hypothetical muon with transverse momentum \( p_T^{\mu} \) and weights the event by \( F(p_T^h, p_T^{\mu}) \). All other event parameters (such as, for instance, the missing energy) are updated to account for this muon. Each event can be used many times for different \( p_T^{\mu} \), as long as \( F(p_T^h, p_T^{\mu}) \neq 0 \).

As an example, we give the expression for the transverse mass distribution

\[
\frac{d\sigma^\mu}{dm_T} = \int dp_T^h dp_T^{\mu} F(p_T^h, p_T^{\mu}) \frac{d\sigma^h}{dp_T^h} \delta(m_T - \sqrt{2 p_T^{\mu} p_T^{\nu} (1 - \cos\phi_{\mu\nu})}).
\]

(7.5.6)

The transverse momentum of the neutrino, \( p_T^{\nu} \), and the angle \( \phi_{\mu\nu} \) between the muon and the neutrino both depend on \( p_T^{\mu} \).

Our aim to create an event sample which is representative of real decay background events thus reduces to the calculation of the probability density function for pion and kaon decays. Once this is known, we can use real data (inclusive hadron events) to generate all kinds of distributions to be subsequently subjected to the W/Z selection criteria.

In the following, we first compute the decay probabilities then discuss the data we use and finally present background estimates.
7. Data Analysis

7.5.2 Calculation of the Probability Density Function for Decays: \( F(p_T^h, p_T^\mu) \)

A calculation of the decay probability based on the free decay length alone is not sufficient. It neglects, for instance, that not all hadrons pass the muon selection criteria. Moreover, the reconstructed \( p_T \) of the CD track is generally different from the real \( p_T \) of the decaying hadron, which is partially due to the finite resolution of the CD, but mainly because of the kink-like structure of the track (see also section 7.2.3). Therefore, a Monte Carlo simulation of pions and kaons decaying in the UA1 detector was performed. This naturally takes into account the effect of track reconstruction and allows to apply the same muon selection criteria, such as track quality and matching requirements, as has been used for the data. The result of the simulation is a probability density \( F(p_T^h, p_T^\mu) \) which yields for each possible hadron (pion or kaon) transverse momentum the probability to reconstruct a muon passing the selection criteria, as a function of the measured muon transverse momentum.

In more detail, the procedure works as follows. Pions and kaons were generated, one particle per event, flat in azimuth, flat in pseudorapidity and flat in \( p_T \) from 3 to 30 GeV (in bins of 1 GeV). Each particle was allowed to decay. If, after starting at a nominal vertex, no decay occurred before the hadron calorimeter, the event was rejected. For the remaining events, a simulation of the calorimetry and the muon chambers was performed. Events with no hits in the muon chambers were rejected. The entire event was reconstructed for the rest.

For the CD reconstruction, the vertex was forced to be at the position of the true Monte Carlo vertex. This is consistent with the treatment of muons in real data. There, the muon track is excluded from the fit of the event vertex (section 7.3.4), so that the latter can be added as an independently measured point.

The sample of reconstructed events was subsequently passed through the inclusive muon selection program (apart from the cut on the muon transverse momentum) discussed in section 7.3. It also had to pass the more stringent cuts on the track quality specific to the W and Z selection. The probability density functions \( F_\pi(p_T^h, p_T^\mu) \) and \( F_K(p_T^h, p_T^\mu) \) were constructed as tables in bins of \( 1 \times 1 \text{ GeV}^2 \), for which the normalization, \( S_h(p_T^h, p_T^\mu) \), for each bin is given by

\[
S_h(p_T^h, p_T^\mu) = \frac{N(p_T^h, p_T^\mu)}{N_g(p_T^h)} \times T. \tag{7.5.7}
\]

In this equation, \( h \) stands for pion or kaon, \( N(p_T^h, p_T^\mu) \) is the number of reconstructed muons with transverse momentum in bin \( p_T^h \) originating from pions/kaons with transverse momentum in bin \( p_T^h \), \( N_g(p_T^h) \) is the number of generated pions/kaons with transverse momentum in bin \( p_T^h \) and \( T \) is the trigger geometrical acceptance. The latter is calculated for each event with a routine which takes into account different trigger conditions weighted by the integrated luminosity for which they were active.

The fractions of pions and kaons in charged hadrons, which appear in equation (7.5.4), have been measured accurately only at low \( p_T \) by UA2\textsuperscript{118} and UA5\textsuperscript{119}. An extrapolation towards higher \( p_T \) was performed in reference 120, where the fractions were found to be
7.5 Evaluation of the π/K Decay Background

\[
\frac{\pi}{h} = 0.58 \pm 0.03, \quad \text{(7.5.8)}
\]

\[
\frac{K}{h} = 0.23 \pm 0.06, \quad \text{(7.5.9)}
\]

independent of \( p_T \) for \( p_T > 3 \) GeV.

These fractions together with the probability densities \( F_\pi \) and \( F_K \) allow us to construct the final one for hadrons. Figure 7.9 shows \( F(p_T^h, p_T^\mu) \) in the form of decay tables. Each curve gives the probability to reconstruct a muon with \( p_T^\mu \) passing all the selection criteria and originating from a hadron in bin \( p_T^h \). Our final interest only concerns muons with high transverse momenta (\( p_T^\mu > 15 \) GeV). From the figure it is clear, that the probability for a low \( p_T \) hadron to become reconstructed as a high \( p_T \) muon is extremely small. Therefore, a very large number of pion and kaon events had to be generated to reduce the statistical uncertainty on the normalization \( S_h \). For the first and second bin in \( p_T^h \), \( 10^6 \) events/bin were generated, while the others correspond to \( 2 \times 10^5 \) events/bin.

![Decay tables](image)

Figure 7.9: \( F(p_T^h, p_T^\mu) \) in the form of decay tables. A curve labelled \( n \) gives the probability to reconstruct a muon with \( p_T^\mu \) passing all the muon selection criteria and originating from a hadron with \( p_T^h \in [n,n+1] \) GeV.
7.5.3 Selection of an Inclusive Hadron Sample

An inclusive hadron sample was obtained with two different triggers: (i) a Minimum Bias trigger (MB) and (ii) a low $E_T$ jet trigger (LJ). Both cause the experiment to run with about 100% deadtime. The MB trigger consists of a pretrigger only and yields hadrons at the low end of the $p_T$ spectrum. The LJ trigger which requires a calorimeter jet with transverse energy exceeding 10 GeV in addition to a valid pretrigger, is biased towards more energetic events and thus covers the higher part of the inclusive hadron $p_T$ spectrum. The integrated luminosities of the MB and LJ samples are $L_{MB} = 56 \, \mu b^{-1}$ and $L_{LJ} = 0.96 \, nb^{-1}$ respectively.

Starting from these samples, a selection was made (table 7.7) which required events to have a hadron track in the CD of reasonable quality pointing to the muon chambers and with $p_T > 3$ GeV for MB events or $p_T > 5$ GeV for LJ events. For both samples, about 12000 events passed the selection. To create a hadron sample with a continuous $p_T$ spectrum from 3 to 20 GeV, MB data was used in the range from 3 to 11 GeV and LJ data in the range from 11 to 20 GeV, while both subsamples were renormalized to the integrated luminosity of the inclusive muon sample: 4.66 pb$^{-1}$. The renormalization factors for the MB and LJ sample are: $R_{MB} = 4.66 pb^{-1}/56 \mu b^{-1} = 83.214$ and $R_{LJ} = 4.66 pb^{-1}/0.96 nb^{-1} = 4854$ respectively. The new $p_T$ spectrum can then be written as

$$\frac{dN}{dp_T} \equiv \begin{cases} R_{MB} \frac{dN_{MB}}{dp_T} & \text{for } p_T \in [3, 11] \text{ GeV} \\ R_{LJ} \frac{dN_{LJ}}{dp_T} & \text{for } p_T \in [11, 20] \text{ GeV} \end{cases}$$ (7.5.10)

This merged sample together with the decay tables calculated earlier will now be used to compute the decay background in the W and Z samples.

7.5.4 $\pi/K$ Decay Background to $W \to \mu \nu$

As explained in section 7.5.1, each hadron event is treated individually. The hadron is replaced by a hypothetical muon with $p_T^\mu$ in the bin $[n_{\text{min}}, n_{\text{min}}+1]$ GeV, the event parameters are updated to account for the muon and the $W \to \mu \nu$ selection criteria are applied to the event as a whole. The exercise is repeated for the next $p_T^\mu$ bin, and so on, up to $[n_{\text{max}}, n_{\text{max}}+1]$ GeV. Then the next hadron event is considered. The whole procedure is rather time consuming and can be speeded up by splitting the $W \to \mu \nu$ selection into two parts. The first part does not involve the value of $p_T^\mu$ and is applied as a sort of preselection to the hadron sample (table 7.8), while the second is applied each time a $p_T^\mu$ has been chosen.
Table 7.8: W preselection applied to the hadron sample. These cuts are independent of the \( p_T \) assumed for the muon.

- \( f_{xy} < 1.0 \)
- \( f_z < 5.0 \)
- \( I_p(\Delta R=0.4) < 1 \text{ GeV} \)
- no CD jet with \( P_T > 5 \text{ GeV} \) back-to-back with the hadron

For each event that passes the W preselection, the hadron track in the CD is extrapolated up to and through the calorimeters. Those calorimeter cells which are hit are marked. Then the hadron is assigned a \( p_T^h \) (the lowest bin to be considered is [15,16] GeV). The response of the calorimeter to a muon with \( p_T^\mu \) is computed and subtracted from the cells which are marked. Subsequently, jet finding in the calorimeter is redone and the missing transverse energy is recalculated, after which it becomes possible to calculate \( M_T, I_E \) and \( I_J \). If the event passes the W cuts on these three variables, it is weighted by \( F(p_T^h, p_T^\mu) \) and by \( R_{MB} \) or \( R_{LJ} \). The event is reused for the next \( p_T^h \) bins up to [59,60] GeV (for higher values of \( p_T^h \), the weighting factor \( F(p_T^h, p_T^\mu) \) becomes negligible). The sum of the weights yields an absolute prediction for the amount of decay background in the W \( \rightarrow \mu\nu \) sample. It was found to be

\[
\# \pi/K \text{ decays in } W \rightarrow \mu\nu \text{ sample: } 74 \pm 3 \text{ (stat)} + 36 \text{ (syst) events.} \quad (7.5.11)
\]

The statistical error was obtained by adding the weights in quadrature and taking the square root. The systematic error is dominated by the uncertainty in the relative pion and kaon fractions given in equations (7.5.8) and (7.5.9), and the uncertainty in the muon energy subtraction procedure. The former amounts to +25% and -33%, while for the latter only an upper bound of +42% was obtained by redoing the analysis without the cuts on the calorimeter isolation.

Since the inclusive hadron selection for LJ events starts at \( p_T = 5 \text{ GeV} \) (table 7.7), the \( p_T \) interval [5,11] GeV represents a region of overlap for the inclusive MB and LJ samples. This interval was used to check whether MB and LJ events lead to different background predictions. It was found that the predictions agreed within the statistical error.

We finally investigated whether the \( p_T \) interval from 3 to 20 GeV in the inclusive hadron spectrum is sufficiently large. It would be worrying, for instance, if the largest background contribution appears to be coming from hadrons with transverse momenta piling up against one of the bounds of the interval. The \( p_T \) spectrum of hadrons that satisfy the criteria in table 7.8 is shown in figure 7.10 with a scale on the left. The solid dots correspond to MB data renormalized by \( R_{MB} \) and the open circles to LJ data renormalized by \( R_{LJ} \). The error bars show that the MB sample runs out of statistics around \( p_T = 11 \text{ GeV} \), which is why we switch to LJ data at that point. The solid line represents a fit to the data.
Figure 7.10: $p_T$ spectrum of hadrons passing the W preselection (solid dots: MB, open circles: LJ, scale on the left), and probability for a hadron to become reconstructed as a muon with a $p_T > 15$ GeV (triangles, scale on the right).

In the same plot, we show the probability for hadrons to become reconstructed as a muon with $p_T > 15$ GeV with a scale on the right. These points were obtained by integrating $F(p_T^h, p_T^\mu)$ over $p_T^\mu$ between 15 GeV and 60 GeV.

The product of the two spectra (we used the fit, to eliminate the fluctuations occurring at large $p_T$), shown in figure 7.11, gives the $p_T$ spectrum of hadrons that contribute to decay background (before the final cuts on $M_T$, $I_E$ and $I_I$ are made). No pile-up is seen at the boundaries. The bin from 3 to 4 GeV contains no entries, while for $p_T > 20$ GeV we expect at most one or two events. The change in slope around 15 GeV reflects the $p_T^\mu > 15$ GeV cut in the W analysis. The step-like behaviour around 9 GeV, which can also be observed in the probability spectrum in figure 7.10, is not really understood. It must be traced back to the decay tables, but no obvious problems could be found there. Since the decay background will be renormalized in the next section using shape arguments, we will not pay further attention to this specific feature.
7.5 Evaluation of the $\pi/K$ Decay Background

7.5.5 $\pi/K$ Decay Background to $Z \rightarrow \mu^+\mu^-$

The whole analysis described in the previous two subsections was repeated with the cuts adjusted to those of the $Z \rightarrow \mu^+\mu^-$ selection. The background appeared to be small

\[
\#\ \pi/K \text{ decays in } Z \rightarrow \mu^+\mu^- \text{ sample: } < 0.1 \text{ events,} \tag{7.5.12}
\]

and will be neglected from now on.

7.6 Renormalization of the $\pi/K$ Decay Background in the $W \rightarrow \mu\nu$ sample

In the last two sections, we presented the data samples that pass the $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ selections, and computed the contributions from different sources of background. In particular, the background from $\pi/K$ decays was found to be negligible for the $Z$ sample, but dominant for the $W$ sample with a large normalization uncertainty. In table 7.9 we have compiled all background contributions to the $W$ sample from which the number of $W \rightarrow \mu\nu$ events is deduced.

In this section, we exploit the possibility to generate decay background distributions in each kinematical variable of interest. By normalizing signal and background distributions to the one observed in the data we might improve our measurement of the observed number of $W \rightarrow \mu\nu$ events. For this purpose we use the $M_T$ distribution.
Table 7.9: $W \rightarrow \mu \nu$ candidates and backgrounds. If two
errors are present, the first is statistical and the second
systematic.

<table>
<thead>
<tr>
<th>process</th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>Candidates</td>
<td>526 ± 23</td>
</tr>
<tr>
<td>backgrounds:</td>
<td></td>
</tr>
<tr>
<td>$\pi/K$ decays</td>
<td>74 ± 3$^{+36}_{-24}$</td>
</tr>
<tr>
<td>$c\bar{c}/b\bar{b}$ decays</td>
<td>7 ± 4</td>
</tr>
<tr>
<td>$\gamma^* \rightarrow \mu \mu$</td>
<td>16 ± 6</td>
</tr>
<tr>
<td>$Z \rightarrow \mu \mu$</td>
<td>9 ± 1</td>
</tr>
<tr>
<td>$W \rightarrow \mu \nu,\tau \nu$ observed</td>
<td>420 ± 23$^{+25}_{-37}$</td>
</tr>
<tr>
<td>5% $W \rightarrow \tau \nu$</td>
<td>20 ± 2</td>
</tr>
<tr>
<td>$W \rightarrow \mu \nu$ observed</td>
<td>400 ± 23$^{+25}_{-37}$</td>
</tr>
</tbody>
</table>

Figure 7.12 shows the $M_T$ distributions from 30 GeV to 230 GeV for $W$'s (including $W$ decays into $\tau$'s), for heavy flavour decays ($c\bar{c}/b\bar{b}$) and for Drell-Yan processes ($DY/Z^0$). They have been generated with ISAJET and include a full simulation of the UA1 detector. The distribution for $\pi/K$ decays was obtained with the method described in the previous section. Clearly, the shape for $W$'s differs from the background at low $M_T$, while the backgrounds themselves all look very similar. The solid lines represent analytic functions which are fitted to the respective distributions.

The fraction of $W$ events in the data is determined by applying the method of maximum likelihood to the observed $M_T$ distribution. If we define $n_i$ as the number of events in histogram bin $i$, and $N$ as the total number of events in a histogram, then the probability $P_i$ to observe $n_i$(data) events originating from the process $W \rightarrow \mu \nu$ and its respective backgrounds is Poissonian distributed

$$P_i = \frac{\lambda_i^{n_i(data)} \ e^{-\lambda_i}}{n_i(data)!}, \quad (7.6.1)$$

with mean

$$\lambda_i = n_i(W) + n_i(\pi/K) + n_i(c\bar{c}/b\bar{b}) + n_i(DY/Z^0), \quad (7.6.2)$$

where $n_i(W)$ is the number of events in bin $i$ of the corresponding histogram for $W$'s, $n_i(\pi/K)$ is the number of events in bin $i$ of the corresponding histogram for $\pi/K$ decays, and so on. The likelihood $\Lambda$ for the observed distribution is given by the product of probabilities over all bins in the histogram

$$\Lambda = \prod_{i=1}^{\# \text{ bins}} P_i. \quad (7.6.3)$$

The fraction of $W$ events can be obtained by maximizing $\Lambda$, while leaving the $W$ and $\pi/K$
normalizations free. To leave the universality assumption for the decays $W \to \mu\nu$ and $W \to \tau\nu$ in tact, the normalizations for the two processes have to be varied simultaneously. At this stage, one could impose the constraint that the sum of the fractions of the different contributions in equation (7.6.2) adds up to one. However, we believe that the data is not sensitive enough to measure such a constraint. Therefore, we vary the $W$ and $\pi/K$ normalizations independently and write $\lambda_i$ as

$$\lambda_i(\alpha, \beta) = \frac{n_i(W)}{N(W)} \alpha N(\text{data}) + \frac{n_i(\pi/K)}{N(\pi/K)} \beta N(\text{data}) + n_i(\text{c\bar{c}/b\bar{b}}) + n_i(\text{DY/Z}^0),$$

(7.6.4)

where $\alpha$ and $\beta$ are the fraction of $W$ and $\pi/K$ events respectively. For the bin contents $n_i$, we use the smooth functional parametrizations instead of the original distributions given in figure 7.13. After taking the logarithm of $\Lambda$ which changes the product into a sum,

$$\log \Lambda = \sum_{i=1}^{\# \text{ bins}} [n_i(\text{data}) \log(\lambda_i(\alpha, \beta)) - \lambda_i(\alpha, \beta)],$$

(7.6.5)
maximization gives for $\alpha$ and $\beta$: $0.715 \pm 0.050$ and $0.224 \pm 0.040$ respectively, where the errors come from the fit (see also figure 7.13). This corresponds to $376 \pm 26$ W's and $118 \pm 21$ $\pi/K$ decays. Thus the prediction for decay background has increased by about one standard deviation and now agrees well with the observed number of events containing definite and doubtful kinks. Since the shape of the three background distributions is very similar, renormalization of the decay background also takes care of the uncertainties in the other backgrounds. Finally, it should be mentioned that although we varied the W and $\pi/K$ normalizations independently, it is clear from figure 7.13 that the W normalization is almost completely constrained by the observed spectrum for large values of $M_T$, say $M_T > 70$ GeV).

The uncertainty on the normalizations of the $c\bar{c}/b\bar{b}$ and DY/$Z^0$ distributions and on the shape of all background distributions introduces an additional 5% uncertainty on the number of W's observed. The effect of the shape was obtained by varying the parameters in the functional forms of the background distributions (figure 7.12) by one standard deviation. After subtraction of the tau decays, the observed number of $W \rightarrow \mu\nu$ events becomes

$$\# \text{ observed } W \rightarrow \mu\nu \text{ events: } 358 \pm 23 \text{ (stat)} \pm 32 \text{ (syst).} \quad (7.6.6)$$
In figure 7.14 we show that the fitted number of W's and π/K decays also leads to good agreement with Monte Carlo predictions for other distributions. The contribution from decay background alone is indicated by the shaded histograms. The Jacobian peak in the $p_T$ spectrum is washed out by detector resolution and is changed into a falling distribution. The fact that the first bin above threshold contains less entries than the second is merely due to a cut on $M_T$. The decay background has a similar shape, but dies out a little faster with increasing $p_T$. The $p_T$ of the W is measured by the pure missing transverse energy (not corrected for $p_T^E$). It too is dominated by detector resolution, which explains why the shape of the decay background so much resembles the one of the data. The shaded histograms in the two isolation plots show that many hadrons appear to be isolated in the CD inside a cone of $\Delta R = 0.4$, but are still surrounded by some neutral energy flow. Notice the flatness of the $I_F$ distribution for decay background, while the one for the data is clearly falling. However, it is not possible to use this variable to perform a global fit as described above, because of the large systematics involved with the muon energy subtraction procedure in the decay background calculation.

Figure 7.14: a) muon transverse momentum, b) W transverse momentum, c) muon momentum isolation and d) muon energy isolation, for the data (data points), the Monte Carlo including all backgrounds (histograms) and the decay background alone (shaded histograms).
We conclude this section by pointing out that the close resemblance between distributions for \( W \to \mu \nu \) and \( \pi/K \) decays inhibits to select an almost background free \( W \to \mu \nu \) event sample without a severe loss in efficiency. In the next section, a different method is described which attacks this particular problem.

7.7 Selection of \( W \to \mu \nu \) and \( Z \to \mu^+ \mu^- \) Candidates, Method 2

7.7.1 Overall Momentum Fit

For the method described here, we also use as input sample the approximately 4000 events that pass the vertex refit procedure discussed in section 7.3.4. The basic difference with the previous method lies in the definition of the muon momentum. Whereas up to now the momentum measurement only made use of the hits recorded in the CD, the present one uses the full muon detection capability of the UA1 spectrometer. It consists of an Overall Momentum Fit (OMF) of the position information from the drift chambers and the limited streamer tube chambers outside the calorimeter, the fitted (by method one) starting point of the muon track in the CD and the event vertex. The muon transverse momentum defined as such, is indicated by \( p_T^{\mu} \) (OMF). Technical details concerning the fit can be found in a UA1 internal note\textsuperscript{121}.

For high momentum tracks, where multiple scattering in the calorimeter and the iron shielding is small, OMF yields an equally good momentum measurement as the one obtained with the CD. An interesting feature of OMF is its possibility to correct some of the

![Figure 7.15: Difference in the angle \( \phi \) in the xy-plane and dip angle \( \lambda \) of the muon track as reconstructed in the CD and by OMF.](image-url)
CD distorsion of the z-coordinate of a muon track. This is reflected in a disagreement in the measured value of the dip angle $\lambda$ as obtained by the two methods shown in figure 7.15. The same figure also shows that both methods agree well on the track angle $\phi$ in the xy-plane. The reason for the improvement on $\lambda$ is that the resolution in z for the vertex and muon chamber hits is much better than the one obtained from the charge division in the CD, and the long lever arm provided by the muon chambers. However, for the method to work well, it was necessary to undo the shift corrections applied to the muon chambers mentioned in section 7.3.2.

### 7.7.2 $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ Selections

The event selections are similar to the ones listed in tables 7.2 and 7.5, but with $p_T^\mu$ replaced by $p_T^\mu$(OMF) and with two additional requirements: (i) a minimum OMF probability and (ii) consistency between $p_T^\mu$ and $p_T^\mu$(OMF). We will first consider the latter two requirements and discuss their use to discriminate against $\pi/K$ decay background.

All events passing the vertex refit procedure were subject to OMF. Subsequently, the $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ selections were carried out, but with a few relaxed cuts, namely: $p_T^\mu$(OMF) > 12 GeV and $M_T(\mu,\nu) > 20$ GeV for W's or $M(\mu_1,\mu_2) > 20$ GeV for Z's. For these two samples, the fit probability $P(\chi^2)$ of OMF is shown in figure 7.16. It should be flat for good muon tracks; a serious overestimation of the resolution on one of the points that support the fit would lead to an enhancement of the distribution near one, whereas underestimating the resolution would give a pileup of events near zero. Tracks with kinks form another cause of low fit probabilities. Clearly, the enhancement near $P(\chi^2) = 0$ is more pronounced for W than for Z candidates. A cut on this variable might remove some decay background.

![Plot](image.png)

Figure 7.16: Overall Momentum Fit probability for W and Z candidates.
Another indication of background in the sample with $W \rightarrow \mu \nu$ candidates can be obtained by considering the normalized difference in $1/p$, $\Delta(1/p)_N$, between the CD and OMF measurements

$$
\Delta(1/p)_N \equiv \frac{1/p(OMF) - 1/p(CD)}{\sqrt{\sigma_{1/p}^2(OMF) + \sigma_{1/p}^2(CD)}},
$$

(7.7.1)

where $\sigma_{1/p}$ represents the error on the measurement. This difference should be standard normally distributed, since the two methods are nearly independent. Distributions of $\Delta(1/p)_N$ are shown in the upper and lower left hand corner of figure 7.17 for the W and Z sample respectively. The solid lines represent gaussian fits of which the mean and standard deviation have been indicated. The width for Z candidates is only slightly larger than one. However, for W candidates it is more than double in size. It means that the errors in $1/p$ have not been correctly estimated, which typically happens if tracks with kinks contaminate the sample. This will be explained when we discuss figure 7.18 below.

![Figure 7.17: Standardized differences in 1/p between the CD and OMF measurements for W and Z candidates before and after the cuts (see text).](image-url)
The $W \rightarrow \mu \nu$ and $Z \rightarrow \mu^+\mu^-$ selections are completed by the cuts $p_T^{\mu}(OMF) > 15$ GeV and $M_1(\mu,\nu) > 30$ GeV for $W$'s and $M(\mu_1,\mu_2) > 50$ GeV for $Z$'s. In addition, we require a good quality fit: $P(\chi^2) > 0.001$, and good agreement between the two methods: $|\Delta(1/p)_N| < 3$. The normalized differences in $1/p$ after these cuts are shown in the upper and lower right hand corner of figure 7.18. Both histograms now closely resemble a standard normal distribution, thus most background has been rejected.

Although the fit probability and normalized difference in $1/p$ are useful tools to remove $\pi/K$ decay background from the $W$ sample, the most effective cut is: $p_T^{\mu}(OMF) > 15$ GeV. This can be explained by the argument illustrated in figure 7.18. It gives a schematic view of a pion or kaon decay occurring inside the CD (solid line). The reconstructed track, based on CD hits only, with its extrapolation towards the muon chambers is indicated by the dashed line. It crosses the muon chambers far away from the actual hits, because the decay point is not correctly recognized. OMF on the other hand computes the track parameters such, that it gives the best agreement between the segment from the vertex to the track entry in the CD and the segment defined by the hits in the muon chambers. Since these two segments are separated by a large distance and the fit is not constrained in the middle of the trajectory, a set of parameters giving a good fit usually exists. However, the apparent bending of the resulting track is larger than the one obtained from the CD alone, thus pushing the transverse momentum below threshold. The OMF result has been indicated by the dashed-dotted line, and its four supporting points by the crosses. Occasionally, muon hits have also been recorded in the Iarocci chambers residing inside the iron shielding close to the muon chambers. In that case, OMF is extended to include those hits.

It now also is clear why $\Delta(1/p)_N$ can be large for kinks: although a reasonable fit combining the two segments often exists, it is not really representative of the actual track including its kink, yielding a wrong value for $\sigma_{1/p}$.

Having discussed some aspects of OMF, we now summarize the $W/Z$ selection for method 2:

- OMF reconstruction with: $P(\chi^2) > 0.001$ and $|\Delta(1/p)_N| < 3$.
- $W/Z$ selections listed in tables 7.2 and 7.5 with: $p_T^{\mu} \rightarrow p_T^{\mu}(OMF)$.

If, for $Z$ candidates, both muons hit the muon chambers, then both have to pass OMF reconstruction.

In total 305 $W \rightarrow \mu \nu$ candidates and 43 $Z \rightarrow \mu^+\mu^-$ candidates were selected, which were subsequently scanned on the Megatek. All $Z$ candidates were found to be ok, while in the $W$ sample, 270 muons were good, 16 contained a definite kink, while for 19 events it was unclear whether there was a kink or not. Thus, the fraction of kinks (counting only definite ones) dropped from 16% for the previous method to 5% for the present one. In figure 7.19 the $p_T^{\mu}$ spectrum for $W$ candidates selected according to method 1 is compared.
to the \( p_T^\mu \) spectrum for ones selected according to the present method. We find that the biggest loss of events is at low \( p_T^\mu \), which is where most background from \( \pi/K \) decays is expected (see also figure 7.14.a).

7.7.3 Backgrounds to \( W \to \mu \nu \) and \( Z \to \mu^+\mu^- \)

The different background contributions were obtained by scaling the predictions listed in tables 7.4 and 7.6 by the OMF reconstruction efficiency, \( \epsilon_{OMF} \). The contribution from \( \pi/K \) decays was calculated separately. To compute \( \epsilon_{OMF} \) we used a sample of \( W \to \mu \nu \) events generated with ISAJET including a full detector simulation. Applying OMF to the reconstructed events yielded \( \epsilon_{OMF} = 85 \pm 4\% \), where the error is mainly systematic, since the statistical error can be neglected. The efficiency was checked on real data. For this purpose we used \( J/\psi \to \mu^+\mu^- \) and \( Z \to \mu^+\mu^- \) events, which are known to be only slightly contaminated by background. For muons from \( J/\psi \)'s, an additional cut: \( p^\mu > 10 \) GeV was required, to remove low momentum muons which suffer from multiple scattering in the calorimeter and iron absorber. We found \( \epsilon_{OMF} = 87 \pm 4\% \) for the \( J/\psi \) data and \( \epsilon_{OMF} = 87 \pm 3\% \) for the \( Z \) data, consistent with the Monte Carlo result. The error on the \( J/\psi \) result mainly reflects some uncertainty on the background, while the one for \( Z \)'s is of a statistical nature. We use the following value: \( \epsilon_{OMF} = 87 \pm 4\% \).

For the calculation of the decay background, we followed the method outlined in section 7.5. The probability density function for decays, \( F(p_T^h, p_T^\mu) \), was updated to account for OMF reconstruction. It was found that only \( 12 \pm 2\% \) of the decay background predicted for the \( W \) sample in section 7.5.4 survives the OMF cuts.
Table 7.10: $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ candidates and backgrounds using OMF. If two errors are present, the first is statistical and the second systematic.

<table>
<thead>
<tr>
<th>process</th>
<th># events</th>
<th>process</th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>$W \rightarrow \mu\nu$ Candidates backgrounds:</td>
<td></td>
<td>$Z \rightarrow \mu^+\mu^-$ Candidates backgrounds:</td>
<td></td>
</tr>
<tr>
<td>$\pi/K$ decays</td>
<td>9 $\pm$ 4</td>
<td>$\gamma^* \rightarrow \mu^+\mu^-$</td>
<td>1.5 $\pm$ 0.8</td>
</tr>
<tr>
<td>$c\bar{c}/b\bar{b}$ decays</td>
<td>5 $\pm$ 3</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$\gamma^* \rightarrow \mu^+\mu^-$</td>
<td>12 $\pm$ 4</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$Z \rightarrow \mu^+\mu^-$</td>
<td>7 $\pm$ 1</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

$W \rightarrow \mu\nu, \tau\nu$ observed 272 $\pm$ 17 $\pm$ 6 5% $W \rightarrow \tau\nu$ 13 $\pm$ 2

<table>
<thead>
<tr>
<th>process</th>
<th># events</th>
<th>process</th>
<th># events</th>
</tr>
</thead>
<tbody>
<tr>
<td>$W \rightarrow \mu\nu$ observed</td>
<td>259 $\pm$ 17 $\pm$ 7</td>
<td>$Z \rightarrow \mu^+\mu^-$ observed</td>
<td>41.5 $\pm$ 6.6 $\pm$ 0.8</td>
</tr>
</tbody>
</table>

A summary of $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ candidates and backgrounds is given in table 7.10. Although the overall detection efficiency for $W \rightarrow \mu\nu$ events has decreased with respect to method 1, the relative statistical error has improved because of a drop in background from 32% for method 1 to 15% for method 2. Even more important is the removal of $\pi/K$ decay background, which eliminates most of the systematic uncertainty on the observed number of $W \rightarrow \mu\nu$ events. For $Z$ candidates, the second method merely represents a loss of detection efficiency.

![Figure 7.20](image)

Figure 7.20: Transverse mass of muon-neutrino pairs for the data, $W$'s (ISAJET) and background.
The transverse mass of muon-neutrino pairs for W candidates obtained by method 2 is given in figure 7.20. Also shown is the contribution from W's and the sum of all backgrounds listed in table 7.10. The shape can hardly be called Jacobian, but at least it shows a turnover in the first three bins following the expected shape of the distribution for real W's. This means a drastic improvement with respect to the result obtained with method 1 displayed in figure 7.13, which just shows a steeply falling distribution.

Finally, an effort was made to prove that indeed the events in our W sample show a resonance behaviour as expected for a decaying object such as the W boson. This could be achieved by defining the muon transverse momentum as a weighted average of the CD and OMF measurements: $p_T^{\mu}$ and $p_T^{\mu}(OMF)$. Since the measurements are nearly independent, combining the two is allowed. However, the weight of the CD momentum partially annuls the effect of OMF, thereby increasing the background from $\pi/K$ decays. This increase can be compensated by a stronger requirement on the consistency between the two momentum measurements at the expense of $W \rightarrow \mu\nu$ statistics. We required $|\Delta(1/p)_{\eta}| < 1$ in addition to the other W selection criteria and relaxed the cut on the muon transverse momentum to: $p_T^{\mu}(CD&OMF \text{ average}) > 12$ GeV. The $p_T^{\mu}(CD&OMF \text{ average})$ distribution for the 232 events that pass the selection is shown in figure 7.21.b. A Jacobian peak now becomes visible. However, the resolution of our detector is at the limit of displaying this feature. The plot should be contrasted with figure 7.21.a, which shows the $p_T^{\mu}$ distribution for W candidates selected by method 1 with $p_T^{\mu} > 12$ GeV.

### 7.8 Overall Detection Efficiency of the $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ Selections

For the calculation of the geometrical acceptance, the trigger efficiency and the efficiency of the event selections we have used the ISAJET Monte Carlo including a simulation of the UA1 detector. A detailed summary is given in tables 7.11 and 7.12. For the $Z \rightarrow \mu^+\mu^-$ selection, the efficiencies have been split into two categories. One concerns events with only one muon leg of the Z pointing into an active muon chamber area, while for the other both point into active areas.

The coverage in $\varphi$ of the muon chambers as a function of $\eta$ has been given in chapter 5 (figure 5.12). It averages to about 65% for $|\eta| < 1$ and 85% for $1 < |\eta| < 2$. We have two indications for the reconstruction efficiency of muons in the muon chambers. The first is a simple hand calculation based on the measured single drift tube efficiency of 92%. If we define this number to be the hit probability, then the probability to get at least three hits out of four in two projections is given by the binomial expansion: $[(0.92)^4 + 4(0.92)^3(0.08)]^2 = 0.93$. This number neglects additional inefficiencies near edges of chambers and the fact that the bottom chambers have only one projection. Therefore, we prefer to use the result of a Monte Carlo calculation: $87 \pm 1\%$. Since various unknown systematic effects are not simulated in the Monte Carlo, this number should be considered as an upper limit. The combined geometrical acceptance and trigger efficiency was obtained by applying an offline version of the first and second level trigger algorithms to reconstructed $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ events generated with the ISAJET Monte Carlo. It takes into account the different trigger conditions which were active during the run and weights them by luminosity. Furthermore, it contains a detailed bookkeeping of dead areas in the muon chambers. For most of the 1988 run, for example, the bottom chambers were off. The third level trigger efficiency was obtained by running the program offline on real data consisting
7.8 Overall Detection Efficiency of the $W \rightarrow \mu \nu$ and $Z \rightarrow \mu^+ \mu^-$ Selections

![Figure 7.21: Muon $p_T$ spectrum for W candidates a) as measured in the CD and b) from a weighted average of CD and OMF measurements with the requirement $|\Delta(1/p)| < 1$. A Jacobian peak is clearly visible in case b.](image)

of 'standard' or 'non-special' events, of which the special events considered in this thesis form a subsample. The result was: $96 \pm 2\%$. The combined efficiencies are $28 \pm 1\%$ and $57 \pm 3\%$ for W's and Z's respectively.

The effect of the kink and cosmic finders and leakage rejection was studied on various data sets. Their efficiency was found to be $96 \pm 1\%$ per muon, which mainly reflects the inefficiency of the kink finder. For the $Z \rightarrow \mu^+ \mu^-$ subsample with two muons pointing to active areas in the muon chambers, the efficiency is practically 100%, because if one muon is rejected there is always a second to survive. The matching efficiency was evaluated on a sample of cosmics. The quoted uncertainty comes from a chamber to chamber variation.

The efficiency of the vertex refit procedure was calculated with inclusive hadrons from the LJ sample passing the selection criteria listed in tables 7.7 and 7.8 to yield event topologies as close as possible to ones of W/Z events. Each event was subject to a vertex refit twice: once excluding tracks from the fit with $p > 6$ GeV and once with $p > 8$ GeV. Sample statistics did not allow to use a momentum threshold of 10 GeV as was used for the inclusive muon data. However, the threshold dependence of the efficiency is small and is accounted for in the quoted systematic uncertainty. The efficiency was found to decrease
with increasing hadron $p_T$ becoming constant for $p_T > 10$ GeV at a value of $90 \pm 2\%$. Another indication of the refit efficiency was obtained from $Z \rightarrow \mu^+\mu^-$ data, known to be clean. A $Z$ selection was performed on the inclusive muon data with $p_T > 12$ GeV before and after the vertex refit. From this we conclude an efficiency of: $93 \pm 3\%$, consistent with the other result. We chose to give a higher weight to the second result, because of the event topology argument, and use: $92 \pm 3\%$.

At this stage, we have to decide on the definition of the muon momentum. However, since the momentum resolutions of both methods are similar, the efficiency of the $p_T > 15$ GeV and $p_T$ (OMF) > 15 GeV cut in the W/Z selections are practically the same, the former being slightly more efficient. For the $p_T$ cut on the second muon in the Z selection, no difference in efficiency between the two methods was found. The quoted uncertainty was obtained by varying in the Monte Carlo the systematic uncertainty on the momentum resolution, $\sigma$, between 0.5$\sigma$ and 2$\sigma$.

The efficiency of the track quality cuts was also computed with generated W/Z events. However, it is known that our simulation of the CD is optimistic about the quality of tracks. The average number of hits per track, for instance, is higher in the Monte Carlo than what is observed in the data. Two studies were made to correct for these effects, one concerns $f_{xy}$ and $f_z$ and the other the number of hits and projected track length.

For the first correction, distributions of $f_{xy}$ and $f_z$ were made for the data, separately for $W \rightarrow \mu\nu$ and $Z \rightarrow \mu^+\mu^-$ events. Subsequently, the cuts on these variables were expressed as a distance from the mean in units of standard deviations. The cuts on the same distributions in the Monte Carlo were adjusted to reproduce these distances. From the shifts of the cuts in the Monte Carlo the size of the corrections could be deduced. The correction to the combined efficiency of the cuts on $f_{xy}$ and $f_z$ is about 3%.

The discrepancy between data and Monte Carlo in the other two variables is larger. A different method was used here. Events were selected from the inclusive muon data with $p_T > 12$ GeV, if they contained a second track with $p_T > 6$ GeV attached to the muon vertex. This yields a sample of ‘high’ $p_T$ tracks to which no quality cuts have been made yet. The efficiency of the cuts on the number of hits and projected track length for these tracks was compared to those of muons from W and Z decays generated with the Monte Carlo in different bins of $p_T$ and $\eta$. The correction factors (~15%) obtained, showed a relatively small $p_T$ and $\eta$ dependence. A much larger variation of the efficiency correction was found upon introduction of additional $f_{xy}$ and $f_z$ cuts applied to the tracks under consideration. Whereas the efficiencies in the Monte Carlo remained almost constant, the ones in the data strongly changed. This indicates the presence of background in the data, which is partially removed by $f_{xy}$ and $f_z$ cuts. These variations account for most of the quoted systematic uncertainty on the track quality efficiencies.

Both the efficiencies of the isolation cuts and mass cuts were computed by Monte Carlo. The 12% relative uncertainty on the absolute energy scale of the hadron calorimeter has only a small effect on the energy isolation criteria, because these are rather loose. Most of the quoted uncertainty can be attributed to the CD isolation cuts, which depend on the ratio of charged to neutral particles originating from fragmentation processes used in the Monte Carlo. The uncertainty on the mass cuts was obtained by varying the momentum resolution, $\sigma$, between 0.5$\sigma$ and 2$\sigma$. 
Table 7.11: Summary of the $W \rightarrow \mu \nu$ acceptance and efficiencies.

<table>
<thead>
<tr>
<th>Category</th>
<th>$W \rightarrow \mu \nu$</th>
</tr>
</thead>
<tbody>
<tr>
<td>geometrical acceptance and trigger efficiency</td>
<td>$0.28 \pm 0.01$</td>
</tr>
<tr>
<td>kink, cosmic and leakage rejection</td>
<td>$0.96 \pm 0.01$</td>
</tr>
<tr>
<td>matching between CD track and muon chamber track</td>
<td>$0.95 \pm 0.05$</td>
</tr>
<tr>
<td>vertex refit</td>
<td>$0.92 \pm 0.03$</td>
</tr>
<tr>
<td>OMF reconstruction</td>
<td>$0.87 \pm 0.04$</td>
</tr>
<tr>
<td>$p_T^\ell / p_T^{\mu}(OMF) &gt; 15$ GeV</td>
<td>$0.81 / 0.79 \pm 0.02$</td>
</tr>
<tr>
<td>track quality</td>
<td>$0.72 \pm 0.06$</td>
</tr>
<tr>
<td>isolation</td>
<td>$0.83 \pm 0.03$</td>
</tr>
<tr>
<td>$M_T(\mu\nu) / M_T^{\mu(OMF)} &gt; 30$ GeV</td>
<td>$0.98 / 0.95 \pm 0.02$</td>
</tr>
<tr>
<td>overall</td>
<td></td>
</tr>
<tr>
<td>method 1:</td>
<td>$0.111 \pm 0.013$</td>
</tr>
<tr>
<td>method 2:</td>
<td>$0.091 \pm 0.012$</td>
</tr>
</tbody>
</table>
Table 7.12: Summary of the $Z \rightarrow \mu^+\mu^-$ acceptance and efficiencies.

<table>
<thead>
<tr>
<th></th>
<th>$Z \rightarrow \mu^+\mu^-$ one muon in active muon chamber area</th>
<th>$Z \rightarrow \mu^+\mu^-$ both muons in active muon chamber area</th>
</tr>
</thead>
<tbody>
<tr>
<td>geometrical acceptance and trigger efficiency</td>
<td>0.36 ± 0.02</td>
<td>0.21 ± 0.01</td>
</tr>
<tr>
<td>kink, cosmic and leakage rejection</td>
<td>0.96 ± 0.01</td>
<td>1.00 ± 0.01</td>
</tr>
<tr>
<td>matching between CD track and muon chamber track</td>
<td>0.95 ± 0.05</td>
<td>1.00 ± 0.01</td>
</tr>
<tr>
<td>vertex refit</td>
<td>0.92 ± 0.03</td>
<td>0.92 ± 0.03</td>
</tr>
<tr>
<td>OMF reconstruction</td>
<td>0.87 ± 0.04</td>
<td>0.76 ± 0.07</td>
</tr>
<tr>
<td>muon 1: $p_T^\mu / p_T^{\mu \text{(OMF)}} &gt; 15 \text{ GeV}$</td>
<td>0.64 / 0.62 ± 0.02</td>
<td>0.91 / 0.89 ± 0.03</td>
</tr>
<tr>
<td>track quality</td>
<td>0.76 ± 0.06</td>
<td>0.89 ± 0.07</td>
</tr>
<tr>
<td>isolation</td>
<td>0.81 ± 0.02</td>
<td>0.95 ± 0.03</td>
</tr>
<tr>
<td>muon 2: $p_T^\mu / p_T^{\mu \text{(OMF)}} &gt; 10 \text{ GeV}$</td>
<td>0.78 ± 0.02</td>
<td>0.95 / 0.95 ± 0.03</td>
</tr>
<tr>
<td>track quality</td>
<td>0.96 ± 0.02</td>
<td>0.98 ± 0.02</td>
</tr>
<tr>
<td>isolation</td>
<td>0.92 ± 0.01</td>
<td>0.99 ± 0.01</td>
</tr>
<tr>
<td>$M(\mu\mu) &gt; 50 \text{ GeV}$</td>
<td>0.97 ± 0.01</td>
<td>0.98 ± 0.01</td>
</tr>
<tr>
<td>overall</td>
<td>method 1: 0.212 ± 0.025</td>
<td></td>
</tr>
<tr>
<td></td>
<td>method 2: 0.166 ± 0.024</td>
<td></td>
</tr>
</tbody>
</table>
8 Determination of $R \equiv \frac{\sigma_W^{\nu}}{\sigma_Z^{ll}}$ and $\Gamma_W^{\text{tot}}$
8. Measurement of $R \equiv \frac{\sigma^{xy}_{W}}{\sigma_{Z}^{ll}}$ and $\Gamma_{W}^{tot}$
8.1 Introduction

In this final chapter, we present $W^\pm$ and $Z^0$ partial cross sections and their ratio, $R$, derived from UA1 data discussed earlier, and compare the results with other measurements and theoretical predictions. The $W$ total width is inferred from the value of $R$ using relation (2.5.2) and the theoretical input discussed in chapters three and four. At this point, we address again the theoretical uncertainty in the ratio of $W^\pm$ and $Z^0$ total cross sections, $R_\sigma$, and show how it can be reduced. Subsequently, the values of $R$ and $\Gamma^{tot}_W$ are used to put a lower bound on the top quark mass and to exclude a region in ($M_2, \mu, \tan\beta$)-parameter space related to $W^\pm$ and $Z^0$ decays into charginos and neutralinos. We conclude this chapter by discussing what improvements on the experimental determination of $\Gamma^{tot}_W$ can be expected from future measurements.

8.2 $W^\pm$ and $Z^0$ Partial Leptonic Cross Sections

8.2.1 Determination of $\sigma^\mu\nu_W$ and $\sigma^\mu\mu_Z$ from the 1988 and 1989 Data

Experimentally, a cross section for a particular process is related to the observed number of events after background subtraction, the detection efficiency $\varepsilon$ and the integrated luminosity $L$, by

$$\sigma = \frac{N_{obs} - N_{bkg}}{\varepsilon \cdot L}.$$  \hspace{1cm} (8.2.1)

For convenience, results obtained in the previous chapter on the number of $W \rightarrow \mu \nu$ and $Z^0 \rightarrow \mu^+\mu^-$ events observed, the number of background events and the overall detection efficiencies according to both selection methods have been compiled in table 8.1. These results together with an integrated luminosity of: $L(1988/1989) = 4.66 \pm 0.37$ pb$^{-1}$, yield the following cross sections

\begin{align*}
\text{method 1:} & \\
\sigma^\mu\nu_W &= 692 \pm 44 \text{ (stat) } \pm 116 \text{ (syst) pb}, \\
\sigma^\mu\mu_Z &= 58.6 \pm 7.8 \text{ (stat) } \pm 8.4 \text{ (syst) pb}, \\
\text{method 2:} & \\
\sigma^\mu\nu_W &= 609 \pm 41 \text{ (stat) } \pm 95 \text{ (syst) pb}, \\
\sigma^\mu\mu_Z &= 53.6 \pm 8.5 \text{ (stat) } \pm 8.9 \text{ (syst) pb}. \hspace{1cm} (8.2.2.a)
\end{align*}

The statistical errors on these cross sections are purely determined by the observed numbers of events, while contributions to the systematic errors come from the luminosity, the efficiency and the number of background events.

The difference in the results obtained with method 1 and 2 can be explained by systematic uncertainties which are not common to both methods. For the $W$ selection, these are the uncertainty associated with the normalization of the background ($\sim 9\%$) and with the OMF reconstruction ($\sim 5\%$) specific to method 1 and 2 respectively. For the $Z$ selection, the main difference comes from the OMF reconstruction alone ($\sim 7\%$). Since the CD and OMF momentum measurement are independent, the associated systematic uncertainties are
8. Measurement of $R = \sigma_W^{\ell\nu}/\sigma_Z^{\ell\ell}$ and $\Gamma_W^{\text{tot}}$

Table 8.1: A compilation of results discussed in chapter 7 concerning the $W \to \mu\nu$ and $Z^0 \to \mu^+\mu^-$ event selections.

<table>
<thead>
<tr>
<th>process</th>
<th>method</th>
<th>$N_{\text{obs}}$</th>
<th>$N_{\text{bkg}}$</th>
<th>$\varepsilon$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$W \to \mu\nu$</td>
<td>1</td>
<td>526 ± 23</td>
<td>168 ± 32</td>
<td>0.111 ± 0.013</td>
</tr>
<tr>
<td>$W \to \mu\nu$</td>
<td>2</td>
<td>305 ± 17</td>
<td>46 ± 7</td>
<td>0.091 ± 0.012</td>
</tr>
<tr>
<td>$Z^0 \to \mu^+\mu^-$</td>
<td>1</td>
<td>60 ± 7.7</td>
<td>2.0 ± 1.0</td>
<td>0.212 ± 0.025</td>
</tr>
<tr>
<td>$Z^0 \to \mu^+\mu^-$</td>
<td>2</td>
<td>43 ± 6.6</td>
<td>1.5 ± 0.8</td>
<td>0.166 ± 0.024</td>
</tr>
</tbody>
</table>

uncorrelated. Therefore, further small ($\sim 3\%$) differences between the two methods for the $W$ and $Z$ selections can arise from cuts involving the muon momentum ($p_T^\mu$, $M_T(\mu\nu)$, $M(\mu\mu)$). In terms of the uncertainties just mentioned, the total difference between the cross sections obtained with method 1 and 2 amounts to about one standard deviation.

For the $W$ cross section measurement, method 2 means an improvement over method 1, because the background is reduced to an acceptable level, namely from 32% to 15%, without reducing the signal too much. Moreover, the large contribution to the systematic error on the cross section caused by the uncertainty in the size of mainly the $\pi$/$K$ decay background is replaced by a smaller one associated with the OMF reconstruction. Finally, as we have seen in the previous chapter, the $M_T(\mu\nu)$ distribution for events selected with method 2 is better behaved than the one obtained with method 1, that is, it closely resembles the shape expected for real $W$'s.

For the $Z$ cross section measurement, the situation is reverse. Because the event sample selected with method 1 contains hardly any background, nothing is gained by using method 2. In fact, it only leads to a loss in statistics and introduces an additional systematic uncertainty associated with the OMF reconstruction.

In the following, we refer to the result obtained with method 2 for $\sigma_W^{\mu\nu}$, while for $\sigma_Z^{\mu\mu}$ we refer to the other method.

8.2.2 Comparison to Other Measurements and Theoretical Predictions

The only other measurement of the $W^\pm$ and $Z^0$ partial cross sections in the muon channel at $\sqrt{s} = 630$ GeV also comes from the UA1 collaboration and is based on data taken in 1984 and 1985 corresponding to an integrated luminosity of 0.6 pb$^{-1}$. In total, 57 $W \to \mu\nu$ and 15 $Z^0 \to \mu^+\mu^-$ events were observed, for which the backgrounds have been computed to be 3.3 and 0.35 events respectively. The following cross sections were measured

\[
\begin{align*}
\sigma_W^{\mu\nu} &= 580 \pm 80 \text{ (stat) } \pm 80 \text{ (syst) pb,} \\
\sigma_Z^{\mu\mu} &= 61 \pm 17 \text{ (stat) } \pm 6 \text{ (syst) pb,}
\end{align*}
\]

(8.2.3)

which are consistent with our present results. The systematic errors on the old results are smaller than the ones obtained with the new data for several reasons. First of all, the old
8.2 $W^\pm$ and $Z^0$ Partial Leptonic Cross Sections

UA1 apparatus contained an electromagnetic calorimeter in the pseudorapidity range $|\eta| < 3$, allowing a better definition of the neutrino in $W \to \mu \nu$ events and limiting the available space for $\pi/K$ decays to the volume of the CD. Secondly, the Central Driftchamber was operated with a larger gas amplification near the sense wires yielding a better momentum resolution mainly in the charge division coordinate. The good CD resolution actually allowed the observation of a clear Jacobian peak in the $p_T$ spectrum for $W$ events. Lastly, a reference sample of approximately 300 $W \to e\nu$ events collected in the same runs and selected by requiring large $E_T$, isolated electrons made it possible to study systematics effects of the CD track quality cuts applied in the muon analysis, on exactly the same type of events.

Under the assumption of lepton universality for the weak charged and neutral couplings the branching ratios $\text{Br}(W \to \ell\nu\tau)$ and $\text{Br}(Z^0 \to \ell^+\ell^-)$, and hence $\sigma_{W}^{\ell\nu}$ and $\sigma_{Z}^{\ell\ell}$, do not depend on the type of lepton. In that case, our results can be directly compared to measurements in other leptonic channels at the same $\sqrt{s}$. An overview is given in table 8.2.

### Table 8.2: $W^\pm$ and $Z^0$ partial leptonic cross sections from the UA1, UA2, and CDF collaborations.

The first error is statistical and the second systematic. For CDF, the systematic error associated with the luminosity measurement is given separately as a third error. The results obtained in this thesis have been indicated by a $^\dagger$.

<table>
<thead>
<tr>
<th>Source</th>
<th>$\sqrt{s}$ (GeV)</th>
<th>Channel</th>
<th>$\sigma \cdot \text{Br}$ (pb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>UA1</td>
<td>546</td>
<td>$W \to e\nu$</td>
<td>$500 \pm 80 \pm 70$</td>
</tr>
<tr>
<td>UA1</td>
<td>630</td>
<td>$W \to e\nu$</td>
<td>$510 \pm 180 \pm 90$</td>
</tr>
<tr>
<td>UA2</td>
<td></td>
<td>$W \to e\nu$</td>
<td>$610 \pm 100 \pm 70$</td>
</tr>
<tr>
<td>UA1</td>
<td></td>
<td>$W \to \mu\nu$</td>
<td>$510 \pm 180 \pm 90$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$W \to e\nu$</td>
<td>$580 \pm 50 \pm 70$</td>
</tr>
<tr>
<td>UA2</td>
<td></td>
<td>$W \to e\nu$</td>
<td>$660 \pm 15 \pm 37$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$W \to \mu\nu$</td>
<td>$580 \pm 80 \pm 80$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$W \to \mu\nu$</td>
<td>$609 \pm 41 \pm 95$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$W \to \tau\nu$</td>
<td>$580 \pm 130 \pm 90$</td>
</tr>
<tr>
<td>CDF</td>
<td>1800</td>
<td>$W \to e\nu$</td>
<td>$2060 \pm 40 \pm 130 \pm 310$</td>
</tr>
<tr>
<td>UA1</td>
<td>546</td>
<td>$Z^0 \to e^+e^-$</td>
<td>$39^{+33}_{-26} \pm 4$</td>
</tr>
<tr>
<td>UA2</td>
<td></td>
<td>$Z^0 \to e^+e^-$</td>
<td>$116 \pm 39 \pm 11$</td>
</tr>
<tr>
<td>UA1</td>
<td></td>
<td>$Z^0 \to \mu^+\mu^-$</td>
<td>$90^{+78}_{-46} \pm 12$</td>
</tr>
<tr>
<td>UA1</td>
<td>630</td>
<td>$Z^0 \to e^+e^-$</td>
<td>$68.0 \pm 14.0 \pm 7.0$</td>
</tr>
<tr>
<td>UA2</td>
<td></td>
<td>$Z^0 \to e^+e^-$</td>
<td>$70.4 \pm 5.5 \pm 4.0$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$Z^0 \to \mu^+\mu^-$</td>
<td>$61.0 \pm 17.0 \pm 6.0$</td>
</tr>
<tr>
<td>UA1$^\dagger$</td>
<td></td>
<td>$Z^0 \to \mu^+\mu^-$</td>
<td>$58.6 \pm 7.8 \pm 8.4$</td>
</tr>
<tr>
<td>CDF</td>
<td>1800</td>
<td>$Z^0 \to e^+e^-$</td>
<td>$197 \pm 12 \pm 10 \pm 30$</td>
</tr>
</tbody>
</table>
together with measurements at other values of √s. Apart from our present measurements, all other UA1 results in the table are based on data taken in the years up to 1985. Both the UA2 and CDF results correspond to data taken in 1988 and 1989.

At √s = 630 GeV, where the UA1 and UA2 experiments have gathered most statistics, results are in good agreement. Clearly, UA2 has significantly reduced both statistical and systematic errors, the latter being largely dominated by the uncertainty on the luminosity which is about 5%. Also the old UA1 results in the electron channel suffer less from systematics owing to the good energy resolution of the calorimeter for high E_t electrons.

In figure 8.1.a, measurements of σ_W^L at √s = 630 GeV are compared to the theoretical prediction which is shown as a function of the top mass. This dependence comes in through the leptonic branching ratio as: σ_W^L = σ_W × Br(W → ℓν) = σ_W × Γ_W^L / Γ_W^tot(m_t).

The figure shows the Born term and the O(α_S) and O(α_S^2) QCD corrections. For the W total cross section we have used the value from equation (4.3.13) in chapter four. The uncertainties associated with the scale, the structure functions and the W mass have been added in quadrature and are indicated by the shaded band. The UA1 measurements corresponding to our new result in the muon channel and old ones in the electron and tau channel, lie systematically below the theoretical prediction for a heavy top (m_t ≥ M_W − m_b). The difference, however, is at most one standard deviation. None of these UA1 results is accurate enough to test QCD, and for a heavy top all of them could be explained by a Born level calculation only. The UA2 result with its much smaller errors does imply that higher order QCD corrections are required, but cannot probe the second order correction separately yet. It also slightly favours light values for the top mass, but is entirely consistent with a heavy top. If one assumes top to be heavy, then the UA2 result is at the limit of constraining the theoretical prediction from above. Since the width of the shaded band is dominated by the uncertainty associated with the choice of structure functions, the measurement provides the possibility to discriminate against parametrizations of u- and d-quark densities yielding too large W cross sections.

Figure 8.1.b shows a similar comparison for σ_Z^L. The m_t-dependence of the leptonic branching ratio is smaller than in the previous case. The low statistics of the Z samples with respect to the W samples leads to relatively large experimental errors. Therefore, these measurements are not sensitive to the top mass.

In figure 8.2.a, the results of the W cross section measurements given in table 8.2 are compared to the theoretical prediction versus √s. The latter was computed assuming top to be heavy. Apart from a tendency for all measurements to lie slightly below the predicted value, the functional dependence of the cross section on √s is very well reproduced. Also, the agreement between theory and experiment for the √s-dependence of σ_Z^L, shown in figure 8.2.b, is excellent.
Figure 8.1: Partial leptonic cross section for: a) $W^\pm$ production and b) $Z^0$ production in $p\bar{p}$ collisions at $\sqrt{s} = 630$ GeV. The theoretical prediction is shown as a function of the top mass. Born level, $O(\alpha_s)$ and $O(\alpha_s^2)$ contributions to the cross section have been indicated separately. The shaded band reflects the uncertainty on the second order result, for which the individual contributions have been added in quadrature.
Figure 8.2: Partial leptonic cross section for: a) $W^\pm$ production and b) $Z^0$ production in $p\bar{p}$ collisions versus $\sqrt{s}$. 

8. Measurement of $R = \sigma_W^{\ell\nu}/\sigma_{Z}^{\ell\ell}$ and $I_W^{\text{tot}}$
8.3 The Ratio $R$

8.3.1 $R$ from the 1988 and 1989 Data

As already mentioned in chapter two, the ratio $R$ is determined experimentally from the relation

$$ R = \frac{N_W - N_W^{\text{bkg}}}{\epsilon_W} \frac{\epsilon_Z}{N_Z - N_Z^{\text{bkg}}} $$

where the luminosities cancel in the ratio. By using the results given in table 8.1, a value of $R$ can be calculated. However, this would lead to a rather pessimistic estimate of the experimental errors, because many components of $\epsilon_W$ and $\epsilon_Z$ are correlated and actually cancel in the ratio. Therefore, the error propagation was done by carrying out small Monte Carlo experiments in which values of $R$ were generated according to the following method. In each experiment, values for the observed number of $W$ and $Z$ events were generated according to Poisson statistics, by selecting at random a value $x_W$ and $x_Z$ from Poissonian distributions $P(x_W, N_W)$ and $P(x_Z, N_Z)$, where

$$ P(x, N) = \frac{N^x}{x!} e^{-N}. $$

Values for the number of background events were generated by selecting at random a value $x_W^{\text{bkg}}$ and $x_Z^{\text{bkg}}$ from uniform distributions $U(x_W^{\text{bkg}}, N_W^{\text{bkg}}, \delta N_W^{\text{bkg}})$ and $U(x_Z^{\text{bkg}}, N_Z^{\text{bkg}}, \delta N_Z^{\text{bkg}})$, where

$$ U(x, N, \delta N) = \frac{1}{\sqrt{12} \delta N} \quad \text{for} \quad N - \sqrt{3} \delta N \leq x \leq N + \sqrt{3} \delta N, $$

$$ = 0 \quad \text{otherwise}. $$

The individual components of $\epsilon_W$ and $\epsilon_Z$ were also fluctuated according to uniform distributions with means and errors as given in tables 7.10 and 7.11, while keeping track of partial correlations between efficiencies. The vertex refit efficiency, for instance, entirely cancels in the ratio. However, the errors on the efficiencies associated with cuts on the second muon leg in the $Z^0 \rightarrow \mu^+ \mu^-$ selection do not cancel at all, while ones associated with cuts on the first muon leg cancel against those from the $W \rightarrow \mu \nu$ selection only partially. To compute $R$, we have chosen to use method 2 for $W$'s and method 1 for $Z$'s, because the statistical error on the number of $Z$'s observed dominates the error on the final result.

A million Monte Carlo experiments have been carried out to obtain a smooth distribution in $R$, which after normalization to one yields a likelihood curve. This curve assigns to each value of $R$ a probability for its occurrence given the numbers of $W$ and $Z$ events observed, the numbers of background events and the efficiencies as determined in chapter 7. The logarithm of this likelihood has been indicated by the dashed curve in figure 8.3. From this curve, the following value of $R$ was determined

$$ R(1988/1989) = 10.1^{+1.8}_{-1.6}. $$
8. Measurement of $R \equiv \sigma_W^V/\sigma_Z^V$ and $\Gamma_W^{\text{tot}}$

![Graph showing log(likelihood) versus R for the old, new and combined UA1 measurements.]

Figure 8.3: Log(likelihood) versus R for the old, new and combined UA1 measurements.

### 8.3.2 Combination of the Old and New UA1 Measurements of R

A similar analysis has been carried out for earlier UA1 data\textsuperscript{123}. Values of R were determined in both the electron and muon channel at $\sqrt{s} = 546$ GeV and 630 GeV using the Monte Carlo method described above. Neglecting a very small $\sqrt{s}$-dependence of R and assuming lepton universality for the weak charged and neutral couplings, the four likelihood curves were combined to obtain an overall fit, from which R was determined to be

$$R(1983-1985) = 9.1^{+1.7}_{-1.2}. \quad (8.3.5)$$

The logarithm of the corresponding likelihood has been indicated by the dotted curve in figure 8.3. The old and new result agree well within the errors. To benefit from the largest possible statistics, we have combined both measurements. The combined likelihood curve is also shown in figure 8.3, from which R was found to be

$$R(\text{combined}) = 9.5^{+1.1}_{-1.0}. \quad (8.3.6)$$

The 90% C.L. interval for the combined measurement is given by: $8.0 < R < 11.6$.

Before comparing this measurement of R to the UA2 result and the theoretical prediction, we first discuss how the uncertainty on this prediction can be reduced.
8.3 The Ratio $R$

8.3.3 A Constraint on $R_\sigma$ from Measurements of the Structure Function Ratio $F_2^h/F_2^p$

Theoretically, the ratio of $W^\pm$ and $Z^0$ partial leptonic cross sections, $R$, is related to the ratio of total cross sections, $R_\sigma$, by

$$R = R_\sigma \frac{\text{Br}(W \rightarrow \ell^+\ell^-)}{\text{Br}(Z^0 \rightarrow \ell^+\ell^-)}.$$  \hspace{1cm} (8.3.7)

In chapter four, we showed that by using structure functions from Morfin and Tung, $R_\sigma$ equals 3.19 (3.31) at $\sqrt{s} = 630$ GeV (1800 GeV). However, by choosing different parametrizations for the parton densities, $R_\sigma$ can vary by as much as 8% at $\sqrt{s} = 630$ GeV for the examples given in table 4.3. The quantity $R_\sigma$ largely depends on the $u(x)/d(x)$ ratio of quark densities in the region: $x = \sqrt{t_{\text{FB}}}$. Since this ratio is an almost linear function of the ratio of the $F_2$ structure functions, $F_2^p/F_2^h$, measurements of the latter can be used to constrain $R_\sigma$ as will be shown below. We start the discussion by restricting ourselves to $\sqrt{s} = 630$ GeV. Afterwards we comment on the effect of going to higher $pp$ collider energies.

Using relations (4.3.4)—(4.3.7), $R_\sigma$ is given in lowest order by

$$R_\sigma = 4 \cos^2\theta_W \frac{\int_{\tau_W}^1 \frac{dx}{x} \left( P_{q\bar{q}W^+}(x, \frac{\tau_W}{x}) + P_{q\bar{q}W^-}(x, \frac{\tau_W}{x}) \right)}{\int_{\tau_Z}^1 \frac{dx}{x} P_{q\bar{q}Z^0}(x, \frac{\tau_Z}{x})},$$  \hspace{1cm} (8.3.8)

where we have integrated over the $\delta$-function appearing in (4.3.4). The joint probability densities defined in (4.3.6) and (4.3.7) can be simplified if one accounts for large contributions only. In equation (4.3.6) for instance, the terms containing $\sin^2\theta_C$ are Cabibbo suppressed and can be neglected. Furthermore, at CERN energies the dominant contribution to the $W^\pm$ and $Z^0$ cross sections comes from interactions between valence quarks. Therefore, neglecting contributions from the quark sea one gets

$$R_\sigma = 4 \cos^2\theta_W \cos^2\theta_C \frac{\int_{\tau_W}^1 \frac{dx}{x} \left( u_1(x)\bar{d}_2\left(\frac{\tau_W}{x}\right) + d_1(x)\bar{u}_2\left(\frac{\tau_W}{x}\right) \right)}{\int_{\tau_Z}^1 \frac{dx}{x} \left( n_u u_1(x)\bar{u}_2\left(\frac{\tau_Z}{x}\right) + n_d d_1(x)\bar{d}_2\left(\frac{\tau_Z}{x}\right) \right)},$$  \hspace{1cm} (8.3.9)

where $n_u$ and $n_d$ have been defined in (4.3.3). Since the quark densities in the proton and the antiquark densities in the antiproton are similar, we change the notation as follows: $q_1(x) \rightarrow q(x)$ and $\bar{q}_2(x) \rightarrow q(x)$, yielding

$$R_\sigma = 4 \cos^2\theta_W \cos^2\theta_C \frac{\int_{\tau}^1 \frac{dx}{x} \left( u(x)d\left(\frac{\tau}{x}\right) + d(x)u\left(\frac{\tau}{x}\right) \right)}{\int_{\tau}^1 \frac{dx}{x} \left( n_u u(x)u\left(\frac{\tau}{x}\right) + n_d d(x)d\left(\frac{\tau}{x}\right) \right)}.$$

In the last step, we have neglected the small difference between $\tau_W$ and $\tau_Z$ ($\tau_W = 0.016$ and
8. Measurement of $R \equiv \sigma^\ell_v / \sigma^\ell_Z$ and $\Gamma^\text{tot}_W$

![Figure 8.4: Parton-parton fluxes in pPb collisions: a) $\tau u(x) d(\tau/x)$, b) $\tau u(x) u(\tau/x)$ and c) $\tau d(x) d(\tau/x)$, evaluated at $\tau = \tau_W$ (solid) and $\tau = \tau_Z$ (dashed) using $M_W = 80.0$ GeV, $M_Z = 91.172$ GeV and $\sqrt{s} = 630$ GeV.](image)

$\tau_Z = 0.021)$. This is justified in figure 8.4 where parton-parton luminosities evaluated at $\tau = M_W^2/s$ and $\tau = M_Z^2/s$ are compared at $\sqrt{s} = 630$ GeV. The same figure shows that the main difference between the $u\bar{d}$, $u\bar{u}$ and $d\bar{d}$ luminosities versus $x$ is the normalization, while the shape of the curves in all cases is very similar. This allows us to make a final approximation for $R_\sigma$

$$R_\sigma = 4 \cos^2 \theta_W \cos^2 \theta_C \frac{\int_{\tau}^{1} \frac{dx}{x} \frac{d(\frac{x}{x})}{u(\frac{x}{x})}}{\int_{\tau}^{1} \frac{dx}{x} \left(n_u + n_d \left[\frac{d(x)}{u(x)} \left[\frac{d(\frac{x}{x})}{u(\frac{x}{x})}\right]\right]\right)} + (u \leftrightarrow d), \quad (8.3.11)$$

revealing the $u(x)/d(x)$-dependence of $R_\sigma$.

At the SpP$S$, first and second order QCD corrections increase the $W^\pm$ and $Z^0$ total cross sections by about 40%. Since these corrections involve quark-gluon terms, the $u/d$-dependence of $R_\sigma$ could in principle be spoiled. However, in chapter four we noticed that $R_\sigma$ is practically insensitive to higher order corrections, the reason being that these corrections are dominated by the $\delta(1 - \frac{x}{x})$ term appearing in the q$q$ subprocess. This effect mainly changes the overall normalization for both $\sigma_W$ and $\sigma_Z$, but cancels in the ratio. Therefore, approximation (8.3.11) remains valid in NLO QCD.

In the parton model, the inelastic $F_2$ structure functions for nucleons are related to the quark densities by

$$F_2(x) = \sum_i e_i^2 x q_i(x), \quad (8.3.12)$$

where the sum runs over the valence- and sea-quarks in the nucleon, $e_i$ is the quark charge and for the moment we have neglected a dependence of the quark densities on $Q^2$ (four-momentum transfer squared). Thus, the ratio $F_2^n/F_2^p$ of neutron and proton structure functions can be written as

$$\frac{F_{2n}^p(x)}{F_{2p}^p(x)} = \frac{u(x) + 4d(x) + S(x)}{4u(x) + d(x) + S(x)}, \quad (8.3.13)$$
where $S(x)$ represents the sum over all sea contributions. For $x \to 0$ the ratio $F_2^d/F_2^p \to 1$, because of sea dominance. For larger $x$ values, the ratio can be approximated by

$$\frac{F_2^d(x)}{F_2^p(x)} \approx \frac{d(x)}{u(x)} + \frac{1}{4}, \quad (8.3.14)$$

where the $d(x)$-term in the denominator of (8.3.13) has been neglected, because for the proton $u(x) \approx 2d(x)$.

The ratio $F_2^d/F_2^p$ has been accurately measured by the BCDMS\(^2\) and NMC\(^3\) collaborations by scattering muons off hydrogen and deuterium targets. Measurements have been performed in different bins of $Q^2$, but within the errors no significant $Q^2$-dependence of $F_2^d/F_2^p$ has been observed. This is expected from QCD, which predicts only small differences in the $Q^2$-evolutions of proton and neutron structure functions. Figure 8.5 shows the $x$-dependence of $F_2^d/F_2^p$ as measured by the two collaborations. The values of $F_2^d/F_2^p$ have been averaged over $Q^2$ for each bin of $x$. The error bars on the data points are statistical. The systematic errors have been indicated by the shaded bands. Both data sets are consistent with $F_2^d/F_2^p \to 1/4$ for $x \to 1$, from which follows that $u(x) \gg d(x)$ at large $x$.

The results from the two experiments are in very good agreement and can be parametrized by

$$P(x) = 1 - 1.85x + 2.45x^2 - 2.35x^3 + x^4, \quad (8.3.15)$$

which fulfills $P(0) = 1$, $P(1) = 1/4$ and $dP/dx(1) = 0$. Figure 8.5 shows that for $x < 0.6$, the ratio $F_2^d/F_2^p(x)$ is well described by (8.3.15) with an uncertainty of about $\pm 0.015$. For larger values of $x$, systematic errors become important.

![Figure 8.5: Measurements of $F_2^d/F_2^p(x)$ by the BCDMS\(^2\) and NMC\(^3\) collaborations and the parametrization $P(x)$ defined in (8.3.15) increased and decreased by 0.015.](image)
8. Measurement of $R \equiv \frac{\sigma^v_W}{\sigma^Z}$ and $\Gamma^\text{tot}_W$

Since $R_\sigma$ directly depends on $u/d(x)$ which is linearly related to $F_2^b/F_2^P(x)$ for $x$ not too small, we expect that parametrizations of parton densities which lead to a good description of the measured ratio $F_2^b/F_2^P$ in the relevant $x$-range ($0.1 < x < 0.5$, see figure 8.4) provide the best estimate of $u/d(x)$ and hence of $R_\sigma$. We have computed $F_2^b/F_2^P(x)$ using different sets parametrizations for the quark densities. The results are shown in figure 8.6, where they are compared to $P(x)$ parametrizing the BCDMS and NMC data. On the basis of this figure, we conclude that the EHLQ1 and DO11 structure functions lead to a bad description of $P(x)$, while the others yield a reasonable agreement (the comparison should not be overdone, because several approximations have been made). Computing $R_\sigma$ with the structure functions shown in figure 8.6 except EHLQ1 and DO1 and averaging the results, gives

$$R_\sigma(\sqrt{s} = 630 \text{ GeV}) = 3.24 \pm 0.06. \quad (8.3.16)$$

From the previous discussion, it is clear how this method breaks down when going to higher $p\bar{p}$ CM energies. It relies on the assumption that $W^\pm$ and $Z^0$ bosons are produced mainly by valence quarks and that the sea contribution to $F_2^b/F_2^P$ in the relevant $x$-range is small. However, with increasing $\sqrt{s}$ the parton densities are probed at smaller values of $x$, so that the sea can no longer be neglected. At these higher energies, the dominant uncertainties are the ratio $d(x)/u(x)$ (including sea components) and the charm structure function $c(x)$. It has been proposed to extract $d(x)/u(x)$ at $p\bar{p}$ colliders directly by measuring the asymmetry in the angular distribution of $W$'s as a function of rapidity (the longitudinal boost of the $W$ reflects the difference in hardness of the $u$- and $d$-quark densities). However, the largest sensitivity is attained at forward rapidities, which requires large statistics. In the same
reference it is argued, that \( c(x) \) could be obtained from the ratio \( R \) again measured as a function of the IVB rapidity (this relies on the observation that the rate for \( c \bar{c} \rightarrow Z^0 \) is much smaller than \( c \bar{c} + \bar{c} + \bar{s} \rightarrow W \)). The problem in connection with determining \( \Gamma_w^{\text{tot}} \) is that the measurement of one variable (\( R \)) is not sufficient to determine two unknowns (\( c(x) \) and \( \Gamma_w^{\text{tot}} \)).

The statistics of the present data sample of the CDF collaboration does not allow to measure the quark densities with sufficient accuracy to discriminate among different sets of structure functions. Since the \( F_2^n/F_2 \) method is no longer applicable at \( \sqrt{s} = 1800 \text{ GeV} \), we have determined \( R_\sigma \) from table 4.3 without excluding the value obtained with the DO1 structure function. This gave

\[
R_\sigma(\sqrt{s} = 1800 \text{ GeV}) = 3.27 \pm 0.06. \quad (8.3.17)
\]

8.3.4 Comparison to Other Measurements and Theoretical Predictions

With the above estimate of \( R_\sigma \), we can now compare our measurement of \( R \) to the theoretical prediction using relation (8.3.7). In figure 8.7, \( R \) is shown as a function of the top mass. The shaded band reflects the uncertainty associated with the structure functions and with \( \sin^2 \theta_W \). The change in slope of \( R(m_t) \) around 45 GeV and 75 GeV is associated with the kinematic closing of the decay channel involving top for \( Z^0 \)'s and \( W^\pm \)'s respectively. For comparison, we have also indicated the effect of the number of light neutrino types, \( N_\nu \), on \( R \).

In the same figure, the UA1 and UA2 results are shown together with their 90\% C.L. intervals (for a discussion of the UA2 measurement, see section 2.6.3). The difference in experimental errors on the two measurements of \( R \) is less pronounced than for the

![Figure 8.7: Measurements of R by the UA1 and UA2 collaborations with their 90% C.L. intervals, and the theoretical prediction as a function of the top mass for 2, 3 and 4 light neutrino types.](image)
individual \( W^\pm \) and \( Z^0 \) cross section measurements. This is due to the fact that the present error is dominated by \( Z^0 \) statistics; the UA1 result is based on 112 \( Z^0 \)'s and the UA2 result on 169 \( Z^0 \)'s. The 90% C.L. interval for UA1 does extend, however, to much larger values of \( R \), which is related to the shape of the likelihood distribution shown in figure 8.3.

The two measurements were statistically combined as follows. First a likelihood curve for the UA2 result was created using the Monte Carlo method discussed earlier. Values of \( R \) were generated by using Poisson statistics for the observed numbers of events (2041 \( W^\pm \) and 169 \( Z^0 \) candidates). Each value of \( R \) was fluctuated according to a uniform distribution of which the size was specified by the systematic error on \( R \) (\( \delta R_{\text{syst}} = 0.25 \)). The likelihood curves for the UA1 and UA2 results were subsequently combined, from which the value of \( R \) was found to be

\[
R(\text{UA1 + UA2}) = 9.4 \pm 0.6, \tag{8.3.18}
\]

with a 90% C.L. interval given by: \( 8.4 < R < 10.4 \).

Comparing this value of \( R \) with the curves on the left in figure 8.7 shows that \( N_\nu \) and \( m_t \) are not constrained simultaneously. However, the former has been precisely determined to be three by the LEP experiments. Both measurements favour \( m_\nu \) values smaller than 75 GeV, although neither of them excludes a heavy top at 90% C.L. The combined measurement falls slightly more than one standard deviation below the expectation for a heavy top, but again does not exclude \( m_t > 75 \) GeV at 90% C.L.

### 8.4 The \( W \) Total Width

#### 8.4.1 The UA1 Result

In section 2.5 we discussed the relation between \( \Gamma_W^{\text{tot}} \) and \( R \), which for convenience is repeated here

\[
\Gamma_W^{\text{tot}} = \frac{C}{R}, \quad \text{with: } C = R_\sigma R_\Gamma \Gamma_Z^{\text{tot}}, \tag{8.4.1}
\]

and where we use the values: \( R_\sigma = 3.24 \pm 0.06 \), \( R_\Gamma = 2.672 \pm 0.030 \) and \( \Gamma_Z^{\text{tot}} = 2.498 \pm 0.020 \) GeV. Because the relation between \( \Gamma_W^{\text{tot}} \) and \( R \) is non-linear, it would be incorrect to simply use the most likely value of \( R \) derived above as input to equation (8.4.1). Therefore the error propagation procedure described in section 8.3.1 was repeated, but this time by generating the distribution \( C/R \). Likelihood curves were constructed for the old and new UA1 measurements, which after combination gave the following value for the total width

\[
\Gamma_W^{\text{tot}}(\text{UA1}) = 2.16^{+0.26}_{-0.25} \text{ (exp.)} \pm 0.05 \text{ (th.) GeV}, \tag{8.4.2}
\]

where the additional theoretical uncertainty of 50 MeV comes from the proportionality constant \( C \), which is dominated by structure function uncertainties. This value of \( \Gamma_W^{\text{tot}} \) assumes an unknown top mass which is free to vary between 50 GeV and 200 GeV. In case one presumes top to be light (~ 50 GeV), the value of \( R_\Gamma \) becomes: \( R_\Gamma = 2.690 \pm 0.021 \) (see section 3.5.1), resulting in a 20 MeV increase of \( \Gamma_W^{\text{tot}} \).
8.4 The W Total Width

8.4.2 The Combined UA1, UA2 and CDF Result

Because our value for the proportionality constant C differs substantially from the ones used by UA2 and CDF (see section 2.6), we use their published values of R and re-compute \( \Gamma^\text{tot}_W \) rather than comparing to their published values of \( \Gamma^\text{tot}_W \) directly. Following the strategy of section 8.3.4, we have generated likelihood curves in C/R for both measurements from which values for \( \Gamma^\text{tot}_W \) could be determined (using \( R_0 = 3.27 \pm 0.06 \) for CDF). The following values have been obtained

\[
\Gamma^\text{tot}_W (\text{UA2}) = 2.29 \pm 0.20 \text{ (exp.)} \pm 0.05 \text{ (th.) GeV}, \quad (8.4.3.a)
\]

\[
\Gamma^\text{tot}_W (\text{CDF}) = 2.12 \pm 0.19 \text{ (exp.)} \pm 0.05 \text{ (th.) GeV}. \quad (8.4.3.b)
\]

If one presupposes top to be light, both values should be increased by 20 MeV. Differences between the values of \( \Gamma^\text{tot}_W \) shown here and the ones quoted in section 2.6 can be entirely explained in terms of differences in the assumptions on the input values entering relation (8.4.1). However, we explicitly attribute a 50 MeV theoretical uncertainty associated with the structure function ambiguity, which is not accounted for in the quoted UA2 and CDF results.

The UA1, UA2 and CDF results agree well within the errors. To benefit from the largest possible statistics, we have used the likelihood curves in C/R for the three experiments to obtain the following combined value for the total width

\[
\Gamma^\text{tot}_W (\text{UA1+UA2+CDF}) = 2.20 \pm 0.12 \text{ (exp.)} \pm 0.05 \text{ (th.) GeV}, \quad (8.4.4)
\]

which again should be increased by 20 MeV if one assumes top to be light.

![Graph showing measurements of \( \Gamma^\text{tot}_W \) by the UA1, UA2 and CDF collaborations and the theoretical prediction.](image-url)
8. Measurement of $R \equiv \sigma_{W^+/W^-}^H / \sigma_Z^H$ and $\Gamma_W^{\text{tot}}$

8.4.3 A Lower Limit on the Mass of the Top Quark

The individual experimental results on $\Gamma_W^{\text{tot}}$ (we use the values given in (8.4.2) and (8.4.3) increased by 20 MeV, since the aim is to derive a limit on $m_t$ in the region $m_t \sim 50$ GeV, thus presupposing top to be light) are compared to the theoretical prediction in figure 8.8, which shows the $m_t$-dependence of the total width and the theoretical uncertainty discussed in chapter three. For $m_t > 75$ GeV, the standard model prediction is: $\Gamma_W^{\text{tot}} = 2.06 \pm 0.02$ GeV. The experimental error and theoretical uncertainty have been indicated separately for each data point. All measurements are consistent with the SM prediction for a heavy top, although the UA2 result exceeds this value by about one standard deviation. Of some (small) significance is the fact that the three results lie systematically above the SM prediction without top, which remains true if we decrease the data points by 50 MeV thus choosing a minimum value for $R_\sigma$ allowed by the structure function ambiguity. This can either be interpreted as being caused by some new process giving a small partial contribution $\Delta \Gamma$ to the total width (be it $W \rightarrow t\bar{b}$ or a non-SM decay), or as being caused by some unknown source of systematics which is common to the three $p\bar{p}$ collider experiments.

In order to derive upper limits on the total widths, the likelihood distributions were numerically integrated over the physically allowed region: $\Gamma_W^{\text{tot}} \geq 2.06$ GeV. The following limits have been obtained at 90% C.L.

\begin{align*}
\text{UA1: } \Gamma_W^{\text{tot}} &< 2.62 \text{ GeV}, & (8.4.5.a) \\
\text{UA2: } \Gamma_W^{\text{tot}} &< 2.64 \text{ GeV}, & (8.4.5.b) \\
\text{CDF: } \Gamma_W^{\text{tot}} &< 2.46 \text{ GeV.} & (8.4.5.c)
\end{align*}

From figure 8.8 it is clear, that the bounds on $\Gamma_W^{\text{tot}}$ by UA1 and UA2 hardly constrain the top mass. For UA1, this is due to the relatively large experimental error, while for UA2 it is also caused by the high central value of the measurement. To derive a lower bound on $m_t$ from the CDF result, we first added in quadrature the theoretical uncertainty on the measurement and on the theoretical prediction, yielding: 0.05 GeV. This was subsequently added linearly to the bound in (8.4.5.c) giving: $\Gamma_W^{\text{tot}} < 2.51$ GeV, or

\begin{align*}
\text{CDF: } m_t &> 42 \text{ GeV} & \text{at } 90\% \text{ C.L.} & (8.4.6)
\end{align*}

In figure 8.9, the combined measurement of $\Gamma_W^{\text{tot}}$ is compared to the SM prediction. Because of the reduced experimental error, the limit on $m_t$ can be improved. In this case, the upper bound on the total width was found to be

\begin{align*}
\text{UA1 + UA2 + CDF: } \Gamma_W^{\text{tot}} &< 2.39 \text{ GeV} & \text{at } 90\% \text{ C.L.}, & (8.4.6)
\end{align*}

which after adding linearly a 50 MeV theoretical uncertainty translates into the following lower bound on the top mass

\begin{align*}
\text{UA1 + UA2 + CDF: } m_t &> 49 \text{ GeV} & \text{at } 90\% \text{ C.L.} & (8.4.7)
\end{align*}
Figure 8.9: $\Gamma_w^{\text{tot}}$ as determined from the combined UA1, UA2 and CDF results, the LEP measurement of $\Gamma_Z^{\text{tot}}$ and the theoretical predictions versus the top mass.

Also shown in figure 8.9 is the LEP measurement of $\Gamma_Z^{\text{tot}}$ to illustrate the difference in sensitivity of the two measurements. The relative uncertainty on the LEP measurement is about 0.8%. The relative experimental uncertainty on $\Gamma_w^{\text{tot}}$ is about seven times larger, or 5.5%. However, an additional 2.5% theoretical uncertainty should be added inherent to the method of determining $\Gamma_w^{\text{tot}}$.

8.5 SUSY Limits

We finally present constraints on supersymmetric $W^\pm$ and $Z^0$ decays in the chargino-neutralino sector that follow from the pp collider results discussed above. There are two ways to exclude regions in $(M_2, \mu, \tan\beta)$-parameter space. The first consists of simply subtracting $\Gamma_w^{\text{tot}}(\text{SM without top}) = 2.06$ GeV from the experimental upper bounds on $\Gamma_w^{\text{tot}}$, and then comparing the remaining excess width $\Delta \Gamma$ to the 2-dimensional surfaces as in figure 3.3 a) for different values of $\tan\beta$. However, this method neglects the strong correlation between the $W^\pm$ and $Z^0$ decay widths. As can be seen from figure 3.3 a) and b), large contributions to $\Gamma_w^{\text{tot}}(\text{SUSY})$ also means large contributions to $\Gamma_Z^{\text{tot}}(\text{SUSY})$. Therefore, it is better to use the following method. We interpret the measured values of R as simultaneously constraining $W^\pm$ and $Z^0$ decay widths, by writing R as

$$\frac{R}{R(\text{SM})} = \frac{\Gamma_w^{\text{tot}} / \Gamma_w^{\text{tot}}(\text{SM}) \Gamma_z^{\text{tot}} / \Gamma_z^{\text{tot}}(\text{SM})}{\Gamma_w^{\text{tot}}(\text{SM}) + \Delta \Gamma_w^{\text{tot}}(\text{SM}) \Gamma_z^{\text{tot}}(\text{SM})}.$$ (8.5.1)
8. Measurement of $R \equiv \sigma_{w}^{\ell\ell} / \sigma_{Z}^{\ell\ell}$ and $\Gamma_{w}^{\text{tot}}$

For a SM with a heavy top: $R(\text{SM}) = 10.54 (10.63)$ at $\sqrt{s} = 630$ GeV (1800 GeV). Combining UA1, UA2 and CDF result yields: $R/R(\text{SM}) = 0.913 \pm 0.047$. The corresponding 90% C.L. interval is given by: $0.836 < R/R(\text{SM}) < 0.990$. The interval was extended linearly by 0.017 to account for the structure function uncertainty on $R(\text{SM})$, giving

$$0.819 < \frac{R}{R(\text{SM})} < 1.007 \quad \text{at 90\% C.L.} \quad (8.5.2)$$

We have used this range to exclude domains in the $(M_2, \mu)$-plane for $\tan \beta = 2$ and $\tan \beta = 8$. The results displayed in figure 8.10, show that a large fraction of $(M_2, \mu)$-space kinematically accessible to $W^\pm$ and $Z^0$ decays can be excluded. Although in the central area (corresponding to small positive values of $\mu$) for $\tan \beta = 2$ the partial contributions to the $W^\pm$ and $Z^0$ decay widths become as large as 400 MeV, this region is not covered. This is due to the fact that here $\Delta \Gamma_{W}(\text{SUSY})$ and $\Delta \Gamma_{Z}(\text{SUSY})$ are approximately equal and hence cannot be detected in $R/R(\text{SM})$. This particular region has been excluded by other experiments$^{125}$ based on measurements of $\Gamma_{Z}^{\text{tot}}$.

Figure 8.10: 90% C.L. excluded domains (shaded areas) in the $(M_2, \mu)$-plane for two values of $\tan \beta$ obtained from the combined UA1, UA2 and CDF measurements of $R$. The dashed-dotted contours indicate the limit of the domain kinematically accessible to $W^\pm$ and $Z^0$ decays into charginos and neutralinos.
8.6 Prospects Concerning $\Gamma_W^{tot}$

At present, the determination of $\Gamma_W^{tot}$ using the method discussed above is limited by statistics. Therefore, improvements can be expected if new data becomes available. Recently (in the second half of 1990), the UA2 experiment finished another data taking run in which their W and Z statistics was approximately doubled by a factor two. Unfortunately, this is too small to allow a significant improvement of the present result.

The CDF and DØ experiments at FNAL are scheduled to take data from summer 1991 to 1992. Both expect to gather an integrated luminosity of ~ 25 pb$^{-1}$ after one year of running, and a total of 100 pb$^{-1}$ from runs up to 1995. This will bring down the statistical error on $\Gamma_W^{tot}$ to about 2% to 3% and probably leads to an improved understanding of the detector thus allowing to reduce the systematic error (the present systematic error on the CDF result is 4%). We therefore expect that the individual experimental error on $\Gamma_W^{tot}$ will be of the order of 5% or: 100 MeV, which does not include the structure function uncertainty. Unless better means are found to constrain the choice of structure functions, the additional 50 MeV uncertainty will remain. Therefore, improvements on $\Gamma_W^{tot}$ from future $\bar{p}p$ collider runs will be modest with uncertainties ranging from 100 MeV to 150 MeV per experiment. However, it is clearly more satisfactory to determine $\Gamma_W^{tot}$ in a single experiment than combining three different results.

In 1994, the beam energy at LEP will be increased to more than twice the W mass to allow the study of W production in $e^+e^-$ collisions. At these energies, most W’s are produced in pairs by $e^+e^-$ annihilation via an intermediate photon or $Z^0$ and by neutrino exchange between the electron and positron. A feasibility study concerning measurements of the W mass and width in these types of events has been presented in reference 126. The authors claim that the W width can be obtained from a fit to the two-jet invariant mass spectrum in $e^+e^- \rightarrow W^+W^- \rightarrow \ell\nu jj$ events with a typical uncertainty of 200 MeV.
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55 See Ref. 54 and also M. Consoli et al., in Ref. 51, Vol. 1, p. 7.
56 For the b-quark, \( \rho \) and \( \bar{s}^2_w \) require slight modifications, see the substitution rules in equation (97) of M. Consoli et al.
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99. These are inelastic, non-diffractive events which have been selected by a pretrigger only. The latter is not sensitive to any specific event signatures, hence the name: ‘Minimum Bias’.


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Summary

This thesis describes an experiment performed with the UA1 detector at the SpS at CERN on the production and decay of $W^\pm$ bosons in proton-antiproton collisions. In particular, the total width of the $W^\pm$ resonance has been determined and its implications for the Standard Model and the Minimal Supersymmetric Model have been investigated.

Chapter 1 gives an outline of the historical role of the p$\bar{p}$ collider project at CERN in understanding electroweak and strong interactions between elementary particles. It is pointed out that even today the production of on-shell $W^\pm$ bosons can only be studied at hadron colliders, because the beam energies of existing lepton machines are too low. A direct measurement of the $W^\pm$ total width exists, but with a relative uncertainty of 70%. The one presented in this thesis is based on an indirect method, which allows to bring down this uncertainty to 8%.

The motivation for determining the $W^\pm$ total width is given in chapter 2. It provides a means to probe the mass of the top quark indirectly. For mass values below the kinematic limit, the top quark can give a partial contribution to the $W^\pm$ width via the decay: $W^+ \rightarrow t\bar{b}$ and the charge conjugate decay. Since the top quark has not been observed so far and there are strong reasons to believe that it exists, the topic is of considerable interest. Extensions of the Standard Model such as the Minimal Supersymmetric Model predict additional decay modes for the $W^\pm$ boson involving new elementary particles, for which the associated partial widths can be significant. Our measurement of the $W^\pm$ total width should be sensitive to such contributions. A non-observation of an excess width with respect to the ‘standard’ one translates into bounds on the masses of those particles.

A description of the method used to determine the $W^\pm$ total width sets the stage for the remainder of this thesis. The width is inferred from the ratio R of the $W^\pm$ and $Z^0$ partial leptonic cross sections, the total width of the $Z^0$ boson and some additional theoretical input. The ratio R is derived from UA1 data and forms the subject of chapters 5, 6 and 7. The $Z^0$
Summary

total width is taken directly from LEP, while the theoretical input is discussed in chapters 3 and 4.

In chapter 3 it is shown that the ratio of $W^\pm$ and $Z^0$ partial leptonic widths, required for the determination of the $W^\pm$ total width, is known with a relative uncertainty of 1%. The Standard Model prediction for the $W^\pm$ total width is discussed in full detail. One section is dedicated to a theoretical discussion of $W^\pm$ and $Z^0$ decays into charginos and neutralinos in the Minimal Supersymmetric Model. Chapter 4 gives numerical predictions for the $W^\pm$ and $Z^0$ total cross sections and their ratio, based on next-to-leading order QCD calculations. The observation is made, that the ratio of total cross sections is almost insensitive to higher order QCD corrections. The dominating uncertainty on this ratio comes from the choice of parametrizations for the parton densities in the (anti)proton. It is constrained, however, by requiring consistency of those parametrizations with the observed ratio of $F_2$ structure functions $F_2^p/F_2^n$, as discussed in chapter 8. The ratio of total cross sections, which is also required for the determination of the $W^\pm$ total width, is then known with a relative uncertainty of 2%.

Chapter 5 describes the UA1 apparatus and includes an overview of the trigger and data acquisition system. Due to the removal of the electromagnetic calorimeter from the detector and a gain reduction of the central drift chamber prior to the 1988 and 1989 runs, the apparatus has lost its electron detection capability and has operated with a degraded momentum resolution. After a discussion of the luminosity measurement, event reconstruction and event simulation are treated in chapter 6. Special attention is payed to muon-identification. The distorting effect of the momentum resolution on the Jacobian peak in the $W^\pm$ transverse mass spectrum is illustrated on simulated events. This effect seriously complicates the selection of a clean $W^\pm$ sample.

The data analysis is presented in chapter 7. A selection of $W^\pm$ and $Z^0$ events is carried out, based on the decays: $W^\pm \rightarrow \mu \nu$ and $Z^0 \rightarrow \mu^+\mu^-$. The $Z^0$ sample hardly suffers from background contamination thanks to the clear signature of two high $p_T$ muons. The background in the $W^\pm$ sample, dominated by pions and kaons decaying in flight, is rather large. It has been attacked in two ways. One method uses a Monte Carlo technique to estimate its size. The other employs a different definition of the muon momentum yielding an enhanced rejection power against $\pi/K$ decays. The chapter concludes with a breakdown of the overall $W^\pm$ and $Z^0$ detection efficiencies.

A discussion of the results is given in chapter 8. The $W^\pm$ and $Z^0$ partial leptonic cross sections are compared to theoretical predictions and to existing measurements. We conclude that our results are well described by the QCD improved parton model, but are not sensitive enough to probe individual first and second order QCD corrections. The uncertainty in $R$ is dominated by $Z^0$ statistics, because most systematic errors cancel in the ratio. The measurement agrees well with an earlier one of UA1 and a recent one of UA2. The uncertainty on the inferred value of the $W^\pm$ total width is also dominated by $Z^0$ statistics. To obtain a best possible estimate, UA1, UA2 and CDF results are combined. After comparing this estimate with the Standard Model prediction, we conclude that a top mass lighter than half the $Z^0$ mass is unlikely. Limits on supersymmetric $W^\pm$ decays into charginos and neutralinos are given in terms of the mass matrix parameters. Finally, we do not expect that the uncertainty in the value of the $W^\pm$ total width will be reduced by more than a factor two in the coming decade.
Samenvatting

Dit proefschrift beschrijft een meting uitgevoerd met de UA1 detektor aan het SpP\(S\) op CERN aan de produktie en het verval van W\(\pm\) bosonen in proton-antiproton botsingen. In het bijzonder is de totale breedte van de W\(\pm\) resonantie bepaald en zijn haar consequenties voor het Standaard Model en het Minimaal Supersymmetrische Model onderzocht.

Hoofdstuk 1 schetst de historische rol van het p\(\bar{p}\) collider project op CERN in het begrijpen van elektrozwakke en sterke interacties tussen elementaire deeltjes. Er wordt op gewezen dat zelfs tegenwoordig de produktie van W\(\pm\) bosonen op de massaschil alleen bestudeerd kan worden aan hadron opslagringen, omdat de bundelenergieën van bestaande lepton machines te laag zijn. Er bestaat een direkte meting van de totale breedte van de W\(\pm\), maar met een relatieve onzekerheid van 70%. De meting gepresenteerd in dit proefschrift is gebaseerd op een indirecte methode, die het mogelijk maakt deze terug te brengen tot 8%.

De motivatie voor het bepalen van de totale breedte van de W\(\pm\) wordt gegeven in hoofdstuk 2. Het biedt de mogelijkheid de massa van het top quark indirekt te onderzoeken. Voor massa’s beneden de kinematische limiet, kan het top quark een partiële bijdrage leveren aan de breedte van de W\(\pm\) via het verval: W\(^+\) \rightarrow t\bar{b} en het ladings geconjugeerde verval. Aangezien het top quark nog niet is waargenomen en er sterke redenen zijn om aan te nemen dat het bestaat, is het onderwerp van aanzienlijk belang. Uitbreidingen van het Standaard Model zoals het Minimaal Supersymmetrische Model voorspellen extra vervals kanalen voor het W\(\pm\) boson, waarbij nieuwe elementaire deeltjes betrokken zijn en waarvoor de partiële breedtes groot kunnen zijn. Onze meting van de totale breedte van de W\(\pm\) zou gevoelig moeten zijn voor zulke bijdragen. Het niet observeren van een additionele breedte met betrekking tot de ‘standaard’ breedte vertaalt in limieten op de massa’s van de deeltjes.

Een beschrijving van de gebruikte methode bepaalt de indeling van de rest van het proefschrift. De breedte wordt afgeleid van de verhouding R van de partiëel leptonische werkzame doorsneden van de W\(\pm\) en de Z\(^0\), de totale breedte van de Z\(^0\) en nog enige theoretische input. De verhouding R is bepaald met UA1 data en vormt het onderwerp van de
hoofdstukken 5, 6 en 7. De totale breedte van de $Z^0$ wordt genomen als bepaald bij LEP, terwijl de theoretische input besproken wordt in de hoofdstukken 3 en 4.

In hoofdstuk 3 wordt aangetoond dat de verhouding van de partiële leptonische breedtes van de $W^\pm$ en de $Z^0$, nodig voor het bepalen van de totale breedte van de $W^\pm$, bekend is met een relatieve onzekerheid van 1%. De voorspelling in het Standaard Model voor de totale breedte van de $W^\pm$ wordt in detail besproken. Een paragraaf is gewijd aan een theoretische behandeling van het verval van $W^\pm$ en $Z^0$ bosonen in charginos en neutralinos in het Minimaal Supersymmetrische Model. Hoofdstuk 4 geeft numerieke voorspellingen voor de totale werkzame doorsneden voor $W^\pm$ en $Z^0$ produktie en hun quotiënt, gebaseerd op tweede orde QCD berekeningen. Er wordt opgemerkt, dat het quotiënt van die werkzame doorsneden vrijwel onafhankelijk is van hogere orde QCD correcties. De onzekerheid op het quotiënt wordt gedomineerd door de keuze van parametrisaties voor de parton dichthen in het (anti)proton. Deze wordt beperkt, echter, door te eisen dat de parametrisaties in overeenstemming zijn met de waargenomen verhouding van $F_2$ structuur functies, $F_2^{q}/F_2^{\bar{q}}$, zoals besproken in hoofdstuk 8. Het quotiënt van de werkzame doorsneden, dat ook nodig is voor de bepaling van de totale breedte van de $W^\pm$, is dan bekend met een relatieve onzekerheid van 2%.

Hoofdstuk 5 beschrijft het UA1 apparaat met inbegrip van een overzicht van het trigger en data acquisitie systeem. Ten gevolge van de verwijdering van de elektromagnetische calorimeter en de reductie in de gasversterking van de centrale driftradiator voorafgaande aan de runs in 1988 en 1989, heeft het apparaat de mogelijkheid verloren elektronen te detekteren en heeft het gewerkt met een verslechterde momentum resolutie. Na een bespreking van de luminositeit meting, wordt de reconstructie en simulatie van gebeurtenissen behandeld in hoofdstuk 6. Speciale aandacht wordt gegeven aan muon-identificatie. Het vervormende effect van de momentum resolutie op de Jacobiaanse piek in het transverse massa spectrum van de $W^\pm$ wordt geïllustreerd met behulp van gesimuleerde gebeurtenissen. Dit effect bemoeilijkt een schone selectie van $W^\pm$'s.

De data analyse wordt weergegeven in hoofdstuk 7. Er wordt een selectie van $W^\pm$ en $Z^0$ gebeurtenissen uitgevoerd, die gebaseerd is op de vervalskanalen: $W^\pm \rightarrow \mu\nu$ en $Z^0 \rightarrow \mu^+\mu^-$. De $Z^0$ verzameling heeft weinig last van achtergrond dankzij de duidelijke signatuur van twee muonen met een hoog transversaal momentum. De achtergrond in de $W^\pm$ verzameling, die gedomineerd wordt door vervallende pionen en kaon, is nogal groot. Zij is op twee manieren benaderd. Een methode gebruikt een Monte Carlo techniek om haar grootte te schatten. De ander gebruikt een gewijzigde definitie van het muon momentum, die een beter scheidingsvermogen oplevert ten aanzien van de $\pi/K$ achtergrond. Het hoofdstuk sluit af met een overzicht van de acceptantie en efficiëntie voor $W^\pm$ en $Z^0$ detektie.

De resultaten worden besproken in hoofdstuk 8. De partiële leptonische werkzame doorsneden voor $W^\pm$ en $Z^0$ produktie worden vergeleken met theoretische voorspellingen en met bestaande metingen. We concluderen dat de resultaten goed beschreven worden door het QCD verbeterde parton model, maar niet gevoelig genoeg zijn om individuele eerste en tweede orde correcties te kunnen testen. De onzekerheid in $R$ wordt gedomineerd door de $Z^0$ statistiek, omdat de meeste systematische fouten in het quotiënt eruit delen. De meting is in goede overeenstemming met een eerdere meting van UA1 en een recente van UA2. De onzekerheid in de afgeleide waarde voor de totale breedte van de $W^\pm$ wordt ook gedomineerd door de $Z^0$ statistiek. Om een zo goed mogelijke schatting te verkrijgen, zijn de resultaten van UA1, UA2 en CDF gecombineerd. Uit een vergelijking van deze schatting met
de Standaard Model voorspelling concluderen we, dat het onwaarschijnlijk is dat de massa van het top quark lichter is dan de halve $Z^0$ massa. Limieten op supersymmetrische $W^\pm$ vervallen in charginos en neutralinos worden gegeven in termen van parameters uit de massa matrix. Tot slot, verwachten we niet dat de onzekerheid in de totale breedte van de $W^\pm$ met meer dan een factor twee verkleind zal worden in het komende decennium.
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The author of this thesis was born in The Hague, The Netherlands, on May 6 1961. He received his high school education at the 1st Vrijzinnig Christelijk Lyceum in The Hague. He obtained a Bachelors Degree in Mathematics and Physics in 1982 from the State University of Leiden, and a Masters Degree in Mathematics and Physics in 1987 from the State University of Utrecht. In preparation to his Masters, he worked on the thermal analysis of bismuth-germanate compounds for which he received a Reward from the University of Utrecht, and was a member of the PEP9/TwoGamma experiment at the Stanford Linear Accelerator Centre in California in the U.S. for a period of a year. He received his graduate training from the Dutch Institute for Nuclear and High Energy Physics (NIKHEF) from 1987 to 1991, during which he was based at CERN near Geneva, Switzerland, being a member of the UA1 experiment. Currently, he is employed by the University of Amsterdam as a research associate in the Computer Systems Group of the Faculty of Mathematics and Informatics working on custom VLSI design.