

**Hadronic Cross Section
Measurements
on the Z Resonance
with the L₃ Detector**

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Hadronic Cross Section Measurements on the Z Resonance with the L₃ Detector

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van de Natuurwetenschappen.

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The branch of experimental high energy physics is clearly one which strongly relies on international cooperation. Just as the construction of present-day experiments is hardly conceivable without the joint efforts of many persons, this thesis would have been impossible to write if I had not been able to communicate with people about various subjects.

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Introduction

The basic aim of high energy physics is to understand the existence of the elementary particles found in nature and their interactions. At present all electromagnetic and weak interactions are well described by the so-called Standard Model, proposed in 1967 by S. Glashow, S. Weinberg and A. Salam [1].

Among the outstanding features of the Standard Model (SM) is the unified treatment of the photon and the charged and neutral weak vector bosons $W^\pm$ and Z . The photon is the carrier of electromagnetic forces between the constituents of matter; similarly the weak vector bosons are responsible for the weak forces between these matter constituents. The discovery of the $W^\pm$ and the Z in 1983, by the UA1 and UA2 Collaborations at the CERN $Sp\bar{p}S$ collider, was the experimental confirmation which gave a solid basis to the Standard Model.

The LEP and SLC colliders (at CERN and SLAC respectively) have been built with the intention to measure the Standard Model parameters and to test the validity of the model. At LEP phase I electrons and positrons are collided at center-of-mass energies of approximately 90 GeV. Here the Z particle is expected to be copiously produced. Its existence gives rise to a resonance behaviour in the cross section of final states. In LEP phase II the center-of-mass energy will be raised to approximately twice the $W^\pm$ mass; thus the study of W^+W^- production will become possible. The information coming from both LEP phases is expected to provide stringent tests of the Standard Model.

One of the physics processes occurring in e^+e^- collisions is the production of hadronic final states. The cross section for hadron production has historically been a rich source of valuable information. The reason for this is that hadrons belong to the most copiously produced final states, which makes it possible to measure their cross section with a good relative precision. This has played a role in the determination of the properties of the J/ψ and Υ resonances, and the confirmation of colour as an additional degree of freedom.

Also at LEP phase I energies, the hadronic cross section is large; it should therefore give rise to precise measurements of the most fundamental properties of the Z boson, i.e. its mass and its total decay width.

The layout of this thesis is as follows:
A brief account of the Standard Model theoretical framework for cross section mea-

surements is given in chapter 1. It is demonstrated that some of the most important information derived from these measurements depends, to a large extent, on small corrections. It is therefore of prime importance to understand these model predictions in detail.

Chapter 2 deals with the experimental setup used to perform the measurements. This setup consists of two parts. First the LEP accelerator, needed to produce the e^+e^- collisions, is described. In particular, attention is paid to the calibration of the LEP beam energy, as the uncertainty in this beam energy directly affects the uncertainty in the measured parameters. Second, an overview of the L_3 detector used to record the above collisions is given. In order to obtain interesting information on the Standard Model precise measurements are needed. The apparatus needed to meet this requirement turns out to be quite complex.

The selection of hadronic events is described in chapter 3. The ‘objects’ used in the reconstruction of the data are explained.

The cross sections resulting from this event selection are presented in chapter 4, in conjunction with a study of their systematic uncertainties.

Finally, an analysis of the results in the context of the Standard Model framework is given in chapter 5.

Chapter 1

Theory

1.1 The Standard Model

The Standard Model mentioned in the introduction of this thesis is the culmination of approximately half a century of work by both theoretical and experimental high energy physicists. Many of its features, notably its chiral – parity violating – structure, the existence of a charged weak vector boson and Quantum Electrodynamics (QED) as a symmetry of nature, were either known or at least anticipated for decades. However, the conceptual difficulties involved in the formulation of a unified treatment of electromagnetic and weak forces turned out to be formidable.

Working in the framework of Yang-Mills theories (where interactions between particles are mediated by means of spin-1 vector bosons), weak interactions between the spin- $\frac{1}{2}$ fermions are interpreted as a representation of the product of an $SU(2)$ symmetry (called $SU(2)_W$) and an $U(1)$ symmetry (called $U(1)_Y$). The spontaneous breaking of the $SU(2)_W \otimes U(1)_Y$ symmetry is an important feature of the model. It gives masses to the fermions and (through the so-called ‘Higgs-Goldstone mechanism’) to three out of four intermediate vector bosons. These are the ‘weak’ vector bosons, the $W^\pm$ and the Z particles. In fact, the term ‘weak’ is merely related to the fact that these bosons are heavy: at low momentum transfers for physical processes – such as β -decay – the large vector boson mass appearing in the denominator of the propagator term manifests itself as a weak effective coupling between the four fermions involved. The remaining vector boson, the photon (γ), stays massless in this scheme. In addition, the existence of a neutral scalar particle, the Higgs particle H , is predicted. This particle is the remnant of the $SU(2)$ field responsible for the breaking of the symmetry.

The fermions of the model can be arranged in so-called ‘families’. Table 1.1 classifies the fermions which are presently known (or anticipated), together with the relevant quantum numbers associated with their $SU(2)_W$ and $U(1)_Y$ representations, the third component of the weak isospin (I_3) and the weak hypercharge (Y) respectively. Also shown is the charge Q of each particle (in units of e). These three

quantum numbers are related to each other by the so-called Gell-Mann – Nishijima relation [2]

$$Q = I_3 + \frac{1}{2}Y. \quad (1.1)$$

The factor $\frac{1}{2}$ in the term $\frac{1}{2}Y$ is due to historical reasons. This relation existed before the invention of the Standard Model, when strong interaction isospin was invoked in order to account for the observed regularities in the masses of hadrons. The chiral nature of the Standard Model reflects itself in the fact that the right-handed and left-handed components (denoted by the subscripts R and L respectively) belong to different representations of the symmetry groups.

I_3	Y	Q	Family		
			1	2	3
$\frac{1}{2}$	-1	0	$\begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}_L$	$\begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}_L$
$-\frac{1}{2}$	-1	-1	e_R^-	μ_R^-	τ_R^-
0	-2	-1			
$\frac{1}{2}$	$\frac{1}{3}$	$\frac{2}{3}$	$\begin{pmatrix} u \\ d \end{pmatrix}_L$	$\begin{pmatrix} c \\ s \end{pmatrix}_L$	$\begin{pmatrix} t \\ b \end{pmatrix}_L$
$-\frac{1}{2}$	$\frac{1}{3}$	$-\frac{1}{3}$	u_R	c_R	t_R
0	$\frac{4}{3}$	$\frac{2}{3}$			
0	$-\frac{2}{3}$	$-\frac{1}{3}$	d_R	s_R	b_R

Table 1.1: The leptons and quarks in the Standard Model and their quantum numbers I_3 , Y and Q .

Two remarks should be made here. First, two of these particles, the top quark (t) and the τ neutrino (ν_τ), have not (yet) been observed directly. However, LEP data provide us with indirect information about both of them. An analysis of the total and partial Z decay widths yields information about the number of ‘light’ neutrino species (i.e. neutrinos with $m_\nu < M_Z/2$) and – as will be seen later – the top quark mass influences radiative corrections.

Second, the notion of the above states as being ‘physical’ states is not entirely correct. The quarks indicated above are the *mass* eigenstates of the theory. The most general form of the Standard Model Lagrangian allows interactions (by means of the weak charged currents) between up-type quarks of one family and down-type quarks of another. (It is this feature which is responsible for the weak decays of strange, charmed and beauty particles.) The eigenstates of the *weak interactions* are then obtained by carrying out an SU(3) transformation (there are three families) on – by convention – the down-type quarks. The resulting 3×3 matrix is called the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix.

The masses of fermions and gauge bosons, and the couplings between them, come about essentially because of three interactions.

The first is a natural consequence of the group theoretic structure of the unbroken model. Requiring a local – ‘gauge’ – symmetry this structure leads to the so-called covariant derivative D_μ . Considering only the leptons of the first family and writing the left-handed $SU(2)_W$ doublet and right-handed singlet as¹

$$L_e = \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}_L, \quad R_e = e_R^- \quad (1.2)$$

one obtains

$$D_\mu L_e = (\partial_\mu + \tfrac{1}{2}ig\mathbf{B}_\mu \cdot \boldsymbol{\tau} + \tfrac{1}{2}ig'\mathcal{A}_\mu Y)L_e \quad (1.3)$$

which gives the interaction between the leptons and the gauge bosons $\mathcal{A}_\mu, \mathbf{B}_\mu$. The coupling constants g, g' determine the physical parameters in the broken theory.

The second ingredient is the introduction of the complex scalar Higgs doublet. Its (postulated) Yukawa coupling to the fermions,

$$-g_e[\bar{R}_e(\phi^\dagger L_e) + (\bar{L}_e\phi)R_e], \quad \phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \quad (1.4)$$

will eventually give mass to the fermions.

Lastly, the ‘free-field’ Lagrangian term for the Higgs doublet field gives mass to the gauge bosons. For the quark doublets the Yukawa coupling is extended with a coupling to the right-handed up-type quarks:

$$-g_u[\bar{u}_R(\bar{\phi}^\dagger L_u) + (\bar{L}_u\bar{\phi})u_R] - g_d[\bar{d}_R(\phi^\dagger L_u) + (\bar{L}_u\phi)d_R] \quad (1.5)$$

where $\bar{\phi} = i\tau_2\phi^*$. Note that it is the absence of the up-type coupling for the leptons which makes the neutrinos massless within the framework of the Minimal Standard Model. The right-handed neutrinos may well exist but since they do not couple to any other field they cannot be observed.

After the symmetry breaking and making an appropriate choice of the Higgs potential and the vacuum expectation value of the Higgs doublet field ϕ , one finds that one particular linear combination of the vector boson fields $B_\mu, \mathcal{A}_\mu$ remains massless and is associated with the photon. The orthogonal combination acquires a mass term and becomes the physical Z boson, with a mass

$$M_Z = \sqrt{g^2 + g'^2}v/2 \quad (1.6)$$

where $v/\sqrt{2}$ is the vacuum expectation value of the ϕ^0 field. Similarly, a mass term develops for the fields B_μ^1, B_μ^2 , with the result

$$W_\mu^\pm = \frac{1}{\sqrt{2}}(B_\mu^1 \pm iB_\mu^2), \quad M_W = gv/2. \quad (1.7)$$

¹We follow the notation of reference 3.

The known electric charge unit e is then associated with the coupling of the fermions to the massless photon field A_μ :

$$e = \frac{gg'}{\sqrt{g^2 + g'^2}}. \quad (1.8)$$

Finally, at the tree level (i.e. without any radiative corrections) the so-called ‘weak mixing angle’ θ_W is defined by

$$\tan \theta_W = g'/g \quad (1.9)$$

Phenomenological input is provided by considering the case of muon decay. In the four-fermion theory the muon lifetime is governed by the so-called Fermi constant G_μ . The precisely measured muon lifetime yields a value of $G_\mu = 1.166389(22) \cdot 10^{-5} \text{GeV}^{-2}$. In the context of the Standard Model, as mentioned in the above, this coupling is re-interpreted in terms of the $W^\pm$ mass, yielding

$$M_W^2 = \frac{\pi \alpha}{\sqrt{2} G_\mu \sin^2 \theta_W} \quad (1.10)$$

where α is the well-known fine structure constant. Furthermore, from the above relations it follows that

$$M_Z = M_W / \cos \theta_W. \quad (1.11)$$

The last two relations have been coupled with measurements of neutral current cross sections in low energy neutrino-electron scattering experiments. Interpreted in the context of the Standard Model one obtains an estimate of $\sin^2 \theta_W$. Before LEP startup, the best value obtained this way was $\sin^2 \theta_W = 0.218 \pm 0.011 \pm 0.004$, where the first error denotes the experimental uncertainties and the second one the uncertainties in the theoretical input [4]. The resulting predictions for M_W, M_Z led to the search range for the $W^\pm$ and Z bosons at the $Sp\bar{p}S$ collider.

At tree level the couplings between the fermions and the bosons of the Standard Model look as follows:

$$\begin{aligned}
\gamma \quad & \text{wavy line} \rightarrow \text{fermion line} & = -ieQ_f\gamma_\mu \\
Z \quad & \text{wavy line} \rightarrow \text{fermion line} & = -ie\gamma_\mu(v_f - a_f\gamma_5) \\
W^\pm \quad & \text{wavy line} \rightarrow \text{fermion line} & = -ie\gamma_\mu \frac{(1 - \gamma_5)}{2\sqrt{2}\sin\theta_W} \\
H^0 \quad & \text{dashed line} \rightarrow \text{fermion line} & = -i\frac{m_f}{2M_W\sin\theta_W}
\end{aligned}$$

where the $\gamma_\mu (\mu = 1, \dots, 5)$ are the Dirac γ matrices and the vector and axial-vector coupling constants v_f, a_f have values

$$v_f = \frac{I_3^f - 2Q_f\sin^2\theta_W}{2\sin\theta_W\cos\theta_W}, \quad a_f = \frac{I_3^f}{2\sin\theta_W\cos\theta_W} \quad (1.12)$$

It should be stressed that the relations (1.10- 1.12) are derived in the context of tree level calculations. After radiative corrections, to be discussed in section 1.3, it turns out that these relations are modified.

A notable feature of the model is that the masses of all the fermions (originating from the Yukawa interaction terms in the unbroken theory) are not fixed. Also the mass of the Higgs boson remains a free parameter.

A last remark should be made regarding the group theoretic structure of the Standard Model. Until now only its electroweak structure has been discussed. Another symmetry group yet plays a non-negligible role: all particles in the Standard Model are also representations of the $SU(3)$ group of strong interactions. The leptons, the photon and the weak gauge bosons are one-dimensional representations of this new group as they do not experience the strong interactions; the quarks however are members of the fundamental three-dimensional representation of this $SU(3)$ group and thus come in three ‘colours’. The gauge bosons responsible for the strong interactions are themselves ‘coloured’ as well: they belong to an eight-dimensional representation of this symmetry group. The dynamics related to this symmetry is called Quantum Chromo Dynamics or QCD. Although this thesis is not primarily concerned with QCD, it is nevertheless necessary to properly take into account its consequences. This has two reasons. The first is a theoretical one: the Standard Model is only renormalisable when also the colour degree of freedom is taken into

account. The second reason is an experimental one: the strong interactions do have an effect on the $e^+e^- \rightarrow q\bar{q}$ cross section as will be seen later.

1.2 Lowest Order Cross Sections

The Feynman diagrams representing the three lowest order processes for fermion production in e^+e^- annihilation, $e^+e^- \rightarrow f\bar{f}$, are depicted in figure 1.1 (in the energy region of the Z resonance, the center-of-mass energy is too low to produce W^+W^- pairs).

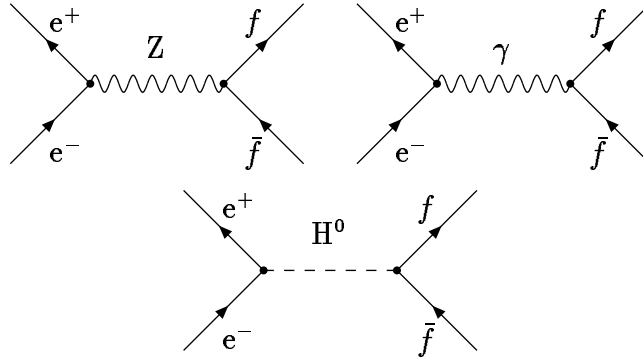


Figure 1.1: *Lowest order graphs for e^+e^- annihilation into fermion pairs $f\bar{f}$ (where $f \neq e^-$).*

Before writing down the lowest order cross sections for unpolarised annihilation it is useful to first make an observation on the order of magnitude of the Higgs exchange contribution. As indicated above, the coupling of the Higgs boson to a fermion is proportional to the mass of this fermion. This means that, considering only the coupling strengths, the contribution of the Higgs channel to the total amplitude for the production of a given fermion f is suppressed (with respect to the γ and Z channels) by a factor $m_e m_f / M_W^2 \alpha$. Even for the heaviest fermion produced, the b quark, this factor is only of the order of $5 \cdot 10^{-5}$. Propagator masses affect this ratio, but the conclusion remains that the Higgs channel gives only a negligible contribution to the total amplitude, even when compared with the photon channel. Consequently, at LEP phase I energies it is always a very good approximation to completely neglect the Higgs channel.

Neglecting the masses of the incoming electron and positron, the lowest order total cross section as a function of the squared center-of-mass energy s becomes

$$\sigma_0^f(s) = \frac{4\pi\alpha^2}{3s} N_C^f \sqrt{1 - 4\mu_f} \left\{ (v_e^2 + a_e^2)(v_f^2(1 + 2\mu_f) + a_f^2(1 - 4\mu_f)) |\chi_0(s)|^2 \right.$$

$$- 2v_e v_f Q_f (1 + 2\mu_f) \text{Re}\chi_0(s) + Q_f^2 (1 + 2\mu_f) \}, \quad (1.13)$$

where

$$\chi_0(s) = \frac{s}{s - M_Z^2 + iM_Z \Gamma_Z}, \quad \mu_f = m_f^2/s \quad (1.14)$$

and Γ_Z is the total Z decay width. The factor N_C^f is introduced in order to take into account the colour degree of freedom. Its value is one in the case of leptons and three for quarks.

The Z decay width is the sum of all the partial widths of the decay of an on-shell Z boson into any of its decay channels. Summing over all fermion species one obtains in lowest order:

$$\Gamma_Z = \sum_f \Gamma_f, \quad \Gamma_f = N_C^f \frac{\alpha}{3} M_Z \sqrt{1 - 4\mu_f} (v_f^2 (1 + 2\mu_f) + a_f^2 (1 - 4\mu_f)). \quad (1.15)$$

Strictly speaking, one might expect that since Γ_Z is formally of higher order, one could neglect it and use the ‘standard’ propagator term for stable particles. Away from the resonance this is a good enough approximation. At LEP phase I energies, however, one measures precisely the resonance curve of the Z boson and with a resolution much better than Γ_Z . This implies that even in lowest order one should use the Z decay width as obtained in eq. 1.15.

In practice, an alternate parametrisation of the weak coupling constants is commonly used:

$$\begin{aligned} g_A^f &= I_3^f = a_f \cdot 2 \sin \theta_W \cos \theta_W, \\ g_V^f &= I_3^f - 2Q_f \sin^2 \theta_W = v_f \cdot 2 \cos \theta_W \sin \theta_W. \end{aligned} \quad (1.16)$$

In terms of these parameters and using eq. 1.10, the expression for the partial widths becomes

$$\Gamma_f = N_C^f \frac{G_\mu M_Z^3}{6\sqrt{2}\pi} \sqrt{1 - 4\mu_f} \{ (g_V^f)^2 (1 + 2\mu_f) + (g_A^f)^2 (1 - 4\mu_f) \}. \quad (1.17)$$

In formula 1.13 one can distinguish three terms. The first one is due entirely to the Z-exchange diagram; it is numerically the most significant term and contains the most interesting physics. It is this channel which leads to the familiar bell-shaped (‘Breit–Wigner’) form of the cross section. The maximum of this function occurs for $s = M_Z^2 (1 + \Gamma_Z^2/M_Z^2)^{1/2}$ and its value then becomes (upon substitution of the above expressions for the partial widths Γ_f)

$$\sigma_{Z,\text{max}}^f = \frac{12\pi \Gamma_e \Gamma_f}{M_Z^2 \Gamma_Z^2} (1 + \Gamma_Z^2/M_Z^2). \quad (1.18)$$

The second and the last term in eq. 1.13 are the (γ, Z) interference and the γ exchange terms respectively. Since at the Z resonance they are numerically small

(of the order of a few percent and a few permille respectively) the qualitative features related to the Z exchange term still dominate.

It should be stressed that the applicability of formula 1.13 does not extend to the Bhabha scattering process $e^+e^- \rightarrow e^+e^-(\gamma)$. The difference arises because the final-state fermions are the same as the initial-state fermions. Consequently, for each of the diagrams depicted in fig. 1.1 there exists a ‘crossed’ diagram, in which the incoming (anti-)fermion line and the outgoing (anti-)fermion line are attached to the same vertex. For these diagrams, using the same arguments as above, the contribution of the Higgs diagram can again be safely neglected. The Z and γ propagators of the two remaining crossed diagrams now involve the momentum transfer $t = (p_{e^-}^{\text{out}} - p_{e^-}^{\text{in}})^2$ rather than the square of the center-of-mass energy s ; for this reason these diagrams are said to constitute the t -channel contribution to the total (lowest-order) scattering amplitude.

Of particular interest in the Bhabha scattering process is the case of small-angle scattering. Using the relation (valid in the center-of-mass system)

$$t = -\frac{s}{2}(1 - \cos \theta) \quad (1.19)$$

the denominator of the t -channel photon propagator is seen to vanish in the limit of forward scattering, giving rise to an infinite total cross section. At small angles, this diagram therefore by far dominates the differential cross section.

The feature of a large cross section in a small part of the solid angle, combined with the fact that it is only very weakly dependent on Standard Model parameters, makes small angle Bhabha scattering very suitable for performing luminosity measurements (see chapter 4); it is generally employed in e^+e^- scattering experiments.

Besides the cross sections there exists another class of measurements which yield important information about Standard Model parameters: the *forward-backward asymmetries*. These asymmetries are defined as

$$A_{FB}^f \equiv \frac{\sigma_F - \sigma_B}{\sigma_F + \sigma_B} \quad (1.20)$$

with

$$\sigma_F = \int_0^1 \frac{d\sigma^f}{d\cos\theta} d\cos\theta, \quad \sigma_B = \int_{-1}^0 \frac{d\sigma^f}{d\cos\theta} d\cos\theta.$$

Evaluating this quantity at $s = M_Z^2$ and neglecting the fermion masses one obtains the following simple relation:

$$A_{FB}^f(M_Z^2) = \frac{3}{4} \mathcal{A}_e \mathcal{A}_f, \quad (1.21)$$

with

$$\mathcal{A}_f = \frac{2v_f a_f}{v_f^2 + a_f^2} = \frac{2g_V^f g_A^f}{(g_V^f)^2 + (g_A^f)^2}. \quad (1.22)$$

The measurement of A_{FB}^f requires the determination of the charge of the final-state fermion. In the case of the process $e^+e^- \rightarrow q\bar{q}$ this implies determining the charge of the quarks responsible for the final state jets (see section 3.2.1). This measurement is not part of the present study. However, in chapter 5 asymmetry measurements performed on leptonic final states will be used in order to increase the precision of the determination of Standard Model parameters.

1.3 Radiative Corrections

The formulas obtained in the previous sections are the results of calculations in the context of lowest order perturbation theory (so-called ‘tree level’ results). As in any perturbative treatment, every amplitude is an infinite series of contributions involving all possible intermediate states. The number of terms of this series that one takes into account is determined by the rate of convergence of the series, the ease of the calculations and by the precision required by the experiment.

A special feature of quantum field theories (as opposed to ‘ordinary’ quantum mechanics) is that the fields have infinitely many degrees of freedom. The consequence is that normally, in the calculation of higher order contributions infinities show up. This apparently cannot be reconciled with the fact that the quantities which are observed in experiments are always finite.

The situation can be saved by the observation that *every* amplitude consists of a perturbative series. This holds for the two-point functions (the propagators) and the vertex functions as well. The higher order corrections applied in these cases imply that the quantities used as input parameters to the theory will in general not be the same as the parameters observed by experiments.

This suggests a re-interpretation of the parameters introduced in the Lagrangian (the ‘bare’ parameters): they are not physical parameters but merely bookkeeping devices. Expressed in terms of the physical parameters any amplitude should assume sensible values. If this holds, order by order in the expansion of the relevant coupling constants and for all amplitudes related to physically observable processes, one is said to deal with a *renormalisable* theory. In such a theory, a finite set of physical observables take over the role of ‘input’ parameters of the theory. Any other observable quantity is then expressed in terms of these input parameters.

The above discussion applies to the Standard Model. This model is a renormalisable theory. (The proof of this assertion in 1972 by G. ’t Hooft and M. Veltman contributed significantly to its acceptance as the ‘standard’ model for describing electroweak interactions [5].) In the present case, the quantities of interest are scattering amplitudes. The LEP collider provides the means to perform precise measurements. The relations between observables and the input parameters are not the same before and after the renormalisation procedure; radiative corrections should therefore be applied.

The experiments commonly used to determine the set of Standard Model input parameters are chosen such that the parameters are measured as directly and independently as possible; in addition they should be measurable with a good precision in order to allow precise tests of the Standard Model. The parameters used (beside the fermion masses) are:

1. the Fermi constant G_μ (from muon decay). G_μ is in fact defined by applying photonic corrections to the μ^- decay process in the framework of the four-fermion theory;
2. the QED coupling constant α (from Thomson scattering, i.e. the low energy limit of Compton scattering);
3. the Z mass M_Z (from the Z lineshape).

Equation 1.10 can be rewritten to relate M_Z , α and G_μ . This relation is however not valid anymore after the application of radiative corrections. The following modification defines the quantity Δr [6]:

$$M_Z^2 \sin^2 \theta_W \cos^2 \theta_W = \frac{\alpha \pi}{\sqrt{2} G_\mu} \frac{1}{1 - \Delta r} \quad (1.23)$$

where $\sin^2 \theta_W$ is defined as $\sin^2 \theta_W = 1 - M_W^2/M_Z^2$ also after renormalisation. Δr is calculable in terms of all fermion and boson masses. It is generally written as

$$\Delta r = \Delta \alpha - \frac{\cos^2 \theta_W}{\sin^2 \theta_W} \Delta \rho + (\Delta r)_{rem} \quad (1.24)$$

where $\Delta \alpha$ represents the change in the ‘running coupling constant’ $\alpha(Q^2)$ from the Thomson scale $Q^2 \approx 0$ to the Z resonance scale $Q^2 = M_Z^2$. Thus one has

$$\alpha(M_Z^2) = \frac{\alpha(0)}{1 - \Delta \alpha}. \quad (1.25)$$

The quantity $\Delta \rho$ directly reflects the radiative corrections to $\sin^2 \theta_W$ and depends on unknown quantities such as m_t and M_H ; the small remaining contribution Δr_{rem} also contains some dependence on ‘new physics’ [6].

The corrections applied to lowest order fermion production can be roughly divided into three categories:

- the *photonic* (QED) corrections; they involve the diagrams where additional photons are either present in the final state (‘real’ corrections) or take part in the intermediate state (‘virtual’ corrections). Since beside the same lowest order scattering these diagrams merely involve well-known QED processes, they are not of particular interest from the point of view of the weak structure of the Standard Model. However, numerically these corrections are very significant. To reliably study other effects they must be taken into account.

- the *electroweak* corrections; they form the main topic of interest. Assuming that the QED corrections applied are reliable, the comparison of lowest order predictions modified by the weak corrections with experimental data provides insight into many interesting aspects of the electroweak structure of the Standard Model.
- the QCD corrections; they are relevant for hadron production only. As in the QED case they comprise both virtual and real corrections. The value of the strong coupling constant α_s is expected to be much larger than the QED coupling constant α . One might therefore expect that the QCD corrections are much harder to account for than the QED ones. However, since the real corrections are relevant for the $q\bar{q}$ final state only, they do not influence the resonance behaviour of the cross section. They can largely be taken into account by means of an overall multiplicative factor.

Each of these three classes of corrections is treated below in somewhat more detail.

1.3.1 Electroweak corrections

Although this division is only unambiguous at one-loop level, the weak corrections can be divided into three types: the propagator corrections, the vertex corrections and the box corrections. A pictorial representation is given in fig. 1.2.

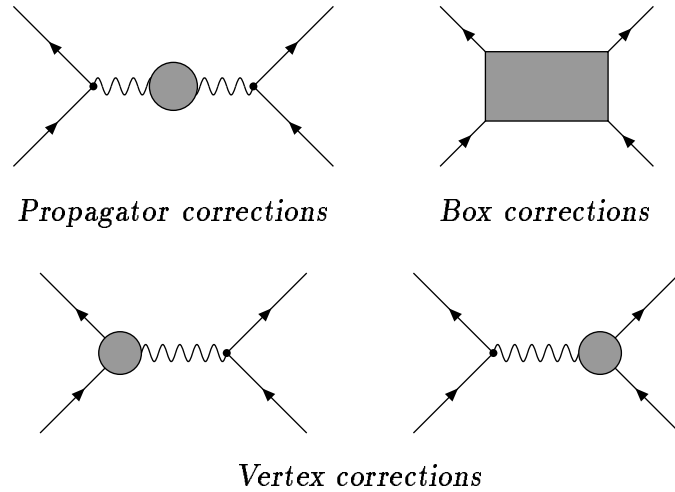


Figure 1.2: *weak corrections to the process $e^+e^- \rightarrow f\bar{f}$.*

Propagator corrections

The propagator corrections are the largest ones contributing to fermion production.

An interesting observation is that, both for intermediate and final states, the states which couple to the Z boson are the same as those which couple to the photon. This leads to the existence of a so-called ‘Z – γ mixing propagator’. By convention, the resulting two mixing terms in the fermion production amplitude are treated as a correction to the Z exchange amplitude.

Consequently, the renormalised (‘dressed’) photon exchange amplitude is almost free from weak corrections. The only remaining effect on this amplitude is the replacement

$$\alpha \longrightarrow \alpha(s) \approx \alpha(M_Z^2). \quad (1.26)$$

In the Z exchange term, electroweak effects cause a so-called ‘self-energy’ correction term in the denominator of the Z propagator. Renormalisation of the Z mass leaves a finite, s -dependent imaginary term which, evaluated at the Z mass, represents the total Z decay width Γ_Z . The s -dependence around the Z resonance is well represented by the following approximation [7]:

$$iM_Z\Gamma_Z \longrightarrow iM_Z\Gamma_Z(s) \approx i\frac{s}{M_Z}\Gamma_Z^{\text{phys}} \quad (1.27)$$

where Γ_Z^{phys} is the physical Z decay width. This is the same quantity as in eq. 1.15, but with weak vertex corrections as well as photonic corrections and QCD corrections (to be discussed below) already applied. This s -dependence of the width is in fact assumed to follow from general quantum field-theoretical arguments [8].

The Z exchange amplitude is subject to additional modifications due to the Z self-energy and Z – γ mixing. This can be approximately accounted for by a re-interpretation of the coupling constants g_A^f, g_V^f :

$$\begin{aligned} g_A^f = I_3^f &\longrightarrow \bar{g}_A^f = \sqrt{\rho}I_3^f, \\ g_V^f = I_3^f - 2Q_f\sin^2\theta_W &\longrightarrow \bar{g}_V^f = \sqrt{\rho}(I_3^f - 2Q_f\sin^2\bar{\theta}_W). \end{aligned} \quad (1.28)$$

This leads to the replacement [9]

$$\frac{\alpha\pi}{\sin^2\theta_W\cos^2\theta_W} \longrightarrow \sqrt{2}G_\mu M_Z^2 \frac{1}{1 - \Delta\rho} \quad (1.29)$$

where $\Delta\rho$ is the same quantity as the one which occurs in eq. 1.24, and $\bar{\theta}_W$ is the so-called *effective* weak mixing angle.

In principle, $\bar{\theta}_W$ and ρ are different for each fermion type. However, in practice *universal* values can be assumed; the only exception is the case of b quarks where due to vertex corrections (see below) $\bar{g}_A^b, \bar{g}_V^b$ are subject to additional modifications. These modifications are taken into account in the Standard Model predictions.

The main contribution to $\Delta\rho$ is given by the term

$$\Delta\rho \approx 3\frac{G_\mu m_t^2}{8\pi^2\sqrt{2}}. \quad (1.30)$$

The relation between $\bar{\theta}_W$ and θ_W is approximately given by

$$\sin^2 \bar{\theta}_W \approx \sin^2 \theta_W - \cos^2 \theta_W \Delta\rho. \quad (1.31)$$

Vertex and box corrections

The vertex corrections originate from diagrams with additional loops in the coupling of the Z or the photon to the external fermions, and from diagrams containing self-energy corrections to the external fermions. Apart from the diagrams involving additional photons (which need to be treated in conjunction with real emission of photons) these corrections are generally small: their effect on the total $f\bar{f}$ cross sections is of the order of 1%.

An exception is the $b\bar{b}$ final state. Since the corrections to $b\bar{b}$ production involve diagrams with a virtual t quark in internal lines these corrections depend on the top mass (quadratically for large m_t). For ‘reasonable’ top mass estimates, $100 \text{ GeV} < m_t < 250 \text{ GeV}$, the corrections range from 1 to 4% [9].

The box corrections are related to diagrams which involve the exchange of two or more gauge bosons. Again, the corrections are separated into diagrams containing additional photons (discussed in the context of photonic corrections) and ‘genuine’ weak corrections, i.e. corrections involving either two Z bosons or two $W^\pm$ bosons. These weak corrections are negligibly small. The reason for this is that their contributions are non-resonant: the Z and W propagators are to be evaluated off-mass-shell.

1.3.2 Photonic corrections

Photonic corrections differ in an essential way from weak corrections because of the mass of the gauge boson involved. At LEP phase I energies, the W and Z boson masses are too heavy to be part of the final state of an event; this implies that all weak corrections are virtual corrections only. On the other hand, the photon is massless so event kinematics allows additional photons in final states (in fact the final states can contain an arbitrary number of photons), i.e. also real (or bremsstrahlung) corrections should be applied. The photonic corrections take into account both classes of radiative corrections.

Virtual and real photonic corrections to fermion production cannot be treated separately. The reason for this is that the contribution from the virtual corrections only (self-energy and vertex corrections) is infinitely negative even after renormalisation. This is due to the exchange of soft ($E_\gamma \rightarrow 0$) photons and is called an infrared (IR) divergence. The differential scattering cross section becomes finite (and positive) only after the addition of the real corrections.

In considering this sum of virtual and real contributions one is helped by the

fact that in the soft limit, the correction can be written as the product of a ‘radiator function’ and a cross section with only weak corrections (this is called factorisation). This property is due to the fact that the interference between initial state radiation and final state radiation is small; fermion production can therefore be viewed as consisting of bremsstrahlung from the incoming e^+ and e^- , followed by the primary annihilation process and bremsstrahlung from the final state fermions.

The bremsstrahlung process involves large terms of the form $\ln(s/m_e^2) \approx 24$ [10]. For this reason the first order corrections are large and higher order corrections cannot be neglected. In general, the ‘soft’ part of the corrections (with a total energy radiated through bremsstrahlung below a suitable cutoff) is *exponentiated*, i.e. corrections to all orders are taken into account. The remaining ‘hard’ photon radiation is treated to second order only.

A convenient way of including the photonic corrections is the so-called ‘structure function’ formalism. Considering only initial state radiation (which gives by far the biggest contribution) one starts out with the weakly corrected cross section and computes the cross section $\sigma(s)$ by means of a convolution of this corrected cross section with radiator kernels describing the bremsstrahlung process. Final state radiation can also be accounted for in this way. The cross section including all corrections can then be expressed as

$$\sigma(s) = \int dx_1 dx_2 dy_1 dy_2 \sigma_0(s') G(x_1, s) G(x_2, s) G(y_1, s'') G(y_2, s'') \quad (1.32)$$

where x_i (y_i) denotes the fraction of longitudinal momentum of the incoming (outgoing) (anti-)fermion after bremsstrahlung, $s' = x_1 x_2 s$ and $s'' = y_1 y_2 s'$; the kernels $G(x, s)$ represent ‘the probability of finding a particle with longitudinal momentum fraction x and virtualness s inside a parent particle’ [11]. The cross section σ_0 comprises not only the weak corrections but also the photonic box corrections and the interference term between the initial and final state radiation terms.

An important consequence of the phenomenon of initial state radiation is the fact that the energy $\sqrt{s'}$ of the primary annihilation process is lowered. This implies that above the resonance, the cross section gets corrections from interactions which take place at the energy of the resonance; it thus develops a so-called *radiative tail*. In the Standard Model, the first order photonic corrections to the cross sections amount to approximately 30%. However, since these photonic corrections are calculable to a good accuracy, the precision of the fully corrected cross sections within the context of the Standard Model is of the order of 0.05% [7].

As mentioned in the previous section, photonic corrections are already included in the definition of the physical decay width Γ_Z^{phys} . These corrections pertain only to final-state photon radiation and are taken into account by an overall multiplicative factor $1 + \delta_{\text{QED}}^f$ where $\delta_{\text{QED}}^f = \frac{3}{4} \frac{\alpha}{\pi} Q_f^2$ is a small term.

The higher order corrections in the treatment of bremsstrahlung not only comprise the cases where there are two or more photons radiated. There are also second

order contributions due to the $f\bar{f}q\bar{q}$ final states. Three higher order processes come into play:

1. radiation of a $q\bar{q}$ pair off an external fermion line in fermion production;
2. radiation of an $f\bar{f}$ pair off a $q\bar{q}$ final state;
3. two-photon processes.

The contributions of the first two of these processes are negligible compared to the photon bremsstrahlung corrections [7]. Only the cross section for the third process becomes large for small invariant masses of the $q\bar{q}$ system. However, in this phase space region the events have a topology rather different from the one resulting from the direct Z/γ annihilation channels. Consequently, these final states are not regarded as being genuine hadronic final states. Instead this process is regarded as background which can be corrected for. The cross section correction following from this process is anyhow small. A more detailed discussion of this process will be given in chapters 3 and 4.

1.3.3 QCD corrections

Like the photonic corrections, the QCD corrections involve the production of additional massless gauge bosons, the gluons. However, unlike the photon the gluon does not couple to leptons. The QCD corrections are therefore applicable to the hadronic final states only. As a consequence, QCD radiative corrections do not affect the energy scale of the ‘hard’ interaction. This greatly simplifies the treatment of the QCD corrections: they can simply be accounted for by an overall multiplicative factor $1 + \delta_{\text{QCD}}^f$ in the total cross section (completely analogous to δ_{QED}^f in the QED case).

Nevertheless, the treatment of QCD corrections remains a nontrivial matter due to other reasons. The running strong coupling constant $\alpha_s(Q^2)$ (≈ 0.12 at LEP energies) is much larger than the QED coupling constant α . The higher order QCD corrections are therefore expected to be significant and hence need to be taken into account. A calculation including the most important terms of $\mathcal{O}(\alpha_s^3)$ has been performed and is presented in reference 12. It should be noted that there are also quark mass effects to be taken into account: the result of these effects is that the partial decay width Γ_b gets a correction which differs slightly from the one applied to the ‘light’ quarks. In all, the QCD corrections amount to an increase of the hadronic decay width by approximately 4%.

Figure 1.3 shows the Standard Model calculation of the total hadronic cross section as a function of $\sqrt{s}$, with and without the corrections discussed above. Evidently, the genuine electroweak corrections change the cross section by a minute

amount only. The effect of the photonic and QCD corrections is to change the peak cross section by approximately 25%. The fermion production cross sections around the Z resonance are nevertheless calculable with a theoretical precision of less than one per mille. This number is obtained from a comparison of several algorithms (see reference 8). This precision enables stringent tests of the validity of the Standard Model and the extraction of physics results (via radiative corrections) which would otherwise be inaccessible at LEP phase I energies.

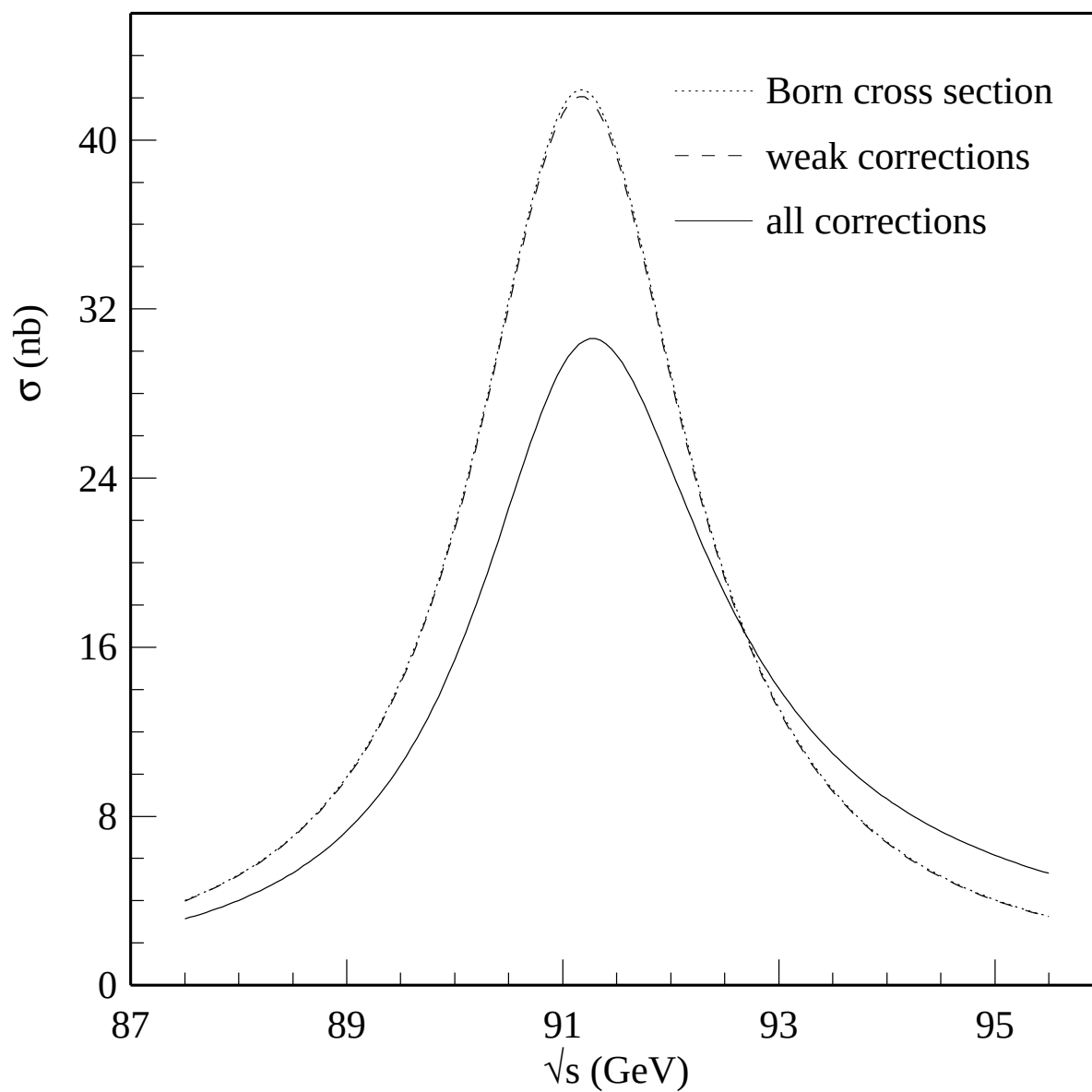


Figure 1.3: *Theoretical prediction for the total hadronic cross section, at tree level and including higher order corrections. Input parameters used are: $M_Z = 91.185 \text{ GeV}$, $m_t = 150 \text{ GeV}$, $M_H = 300 \text{ GeV}$, $\alpha_s = 0.12$.*

Chapter 2

Experimental Setup

2.1 LEP

2.1.1 The LEP accelerator system

The LEP collider, situated at the French-Swiss border near Geneva, currently reaches the highest center-of-mass energies ever achieved in e^+e^- experiments. The design goal for LEP phase I was to cover the Z resonance. Therefore the accelerator had to be able to produce electrons and positrons up to energies of approximately 50 GeV.

The LEP ring has a circumference of approximately 27 km and contains eight straight and eight curved sections. This circumference makes LEP the largest existing storage ring. It was chosen as large as possible (within the constraints imposed by geology and construction costs), in order to reduce the energy loss due to synchrotron radiation suffered by accelerated charged particles. The energy loss per revolution is proportional to E_b^4/ρ where E_b is the particle beam energy and ρ is the radius of curvature [14]. With a circumference of 27 km the energy needed to replenish this energy loss is reduced to a manageable 220 MeV per turn [15].

Each of the eight curved sections contains a set of 31 ‘standard’ magnet cells (of 79 m length) consisting of bending dipole magnets and focusing quadrupole and sextupole magnets. The operation of the ring is organised into so-called ‘fills’, each of which starts out with a period in which the ring is filled with opposing beams of electrons and positrons. These are then accelerated to their required energies and brought into collision in the desired interaction points.

A set of 128 copper radio-frequency (RF) cavities powered by 16 one MW klystrons is used for the acceleration of the beams. The RF cavities are operated at two frequencies, both of which are harmonic numbers of the LEP revolution frequency for ultrarelativistic particles; the difference between the two frequencies is eight times the revolution frequency. The use of two slightly different frequencies causes so-called ‘beating’, a continuous oscillation of the electromagnetic energy between the RF cavities themselves and low-loss storage cavities placed on top of

them. This significantly reduces the power consumption by the acceleration system. Presently, the cavities are gradually being replaced with superconducting ones in order to prepare for the higher energies in the LEP II phase.

Separators (electrostatical plates) are used at the interaction points to precisely steer the electron and positron beams. This serves to maximise the luminosity during physics runs but also to separate the beams (in the transverse plane) during the acceleration phase when the interactions between the beams would destroy them. For geological reasons the plane of the ring is tilted by 1.4° with respect to the horizontal plane. For details on the design of the LEP ring see ref. 16.

For the pre-acceleration of the electrons and positrons LEP uses (as much as possible) already existing facilities. There are five main steps. First electrons and positrons are produced and accelerated to 600 MeV in the LEP Injector Linacs (LIL); these particles are then injected into the Electron Positron Accumulator (EPA) which groups them into bunches. The bunches are subsequently injected into the Proton Synchrotron (PS) and accelerated up to 3.5 GeV. After injection into the Super Proton Synchrotron (SPS) they are accelerated further to 20 GeV. The energy is then high enough for the bunches to be accelerated up to their collision energies in the LEP ring itself.

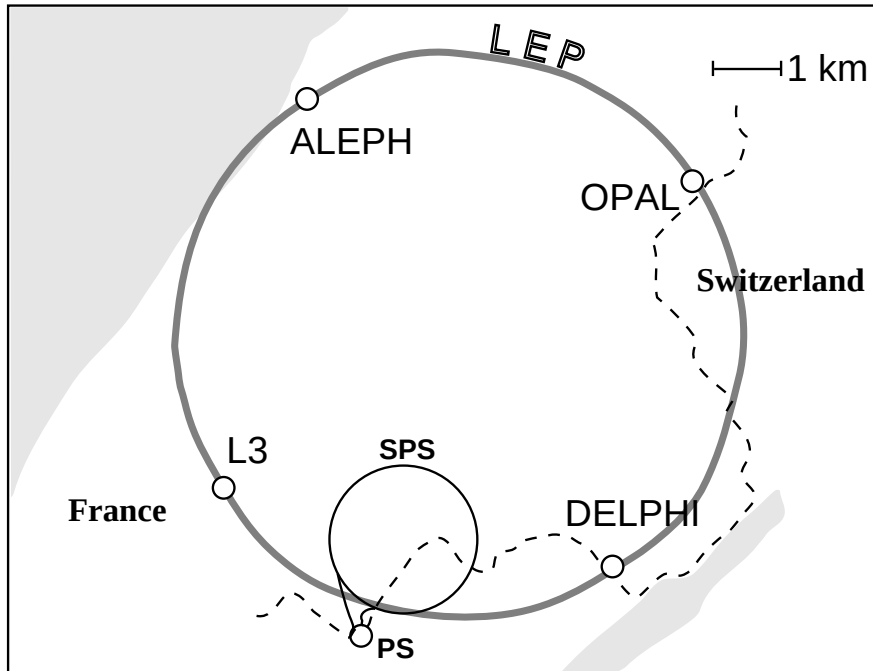


Figure 2.1: *The LEP collider complex.*

From the start-up of LEP in 1989 until the 1992 data taking period, four bunches of electrons were collided with four bunches of positrons. From 1992 onwards a doubling of the number of bunches was effectuated in order to increase the luminosity. A precise timing system ensures that the e^+ and e^- bunches are injected at the time of the RF oscillation ('RF bucket') required for the collisions to actually take place in the experimental interaction points. The oscillations of the particles within a bunch around the ideal ('synchronous') particle trajectory result in a beam energy spread of approximately 50 MeV. The accelerator complex is described in great detail in ref. 17; see also e.g. ref. 18.

The experimental areas are all situated in the straight sections. The even-numbered interaction points (IP) accommodate the LEP detectors L_3 (IP2) [13], ALEPH(IP4) [19], OPAL(IP6) [20] and DELPHI(IP8) [21] respectively. Figure 2.1 presents a general overview of the LEP collider complex.

2.1.2 The LEP energy calibration

Since the measurement of the most elementary properties of the Z, i.e. its mass and decay width, depends directly on the measurement of cross sections as a function of $\sqrt{s}$, any uncertainty in the LEP center-of-mass energies will directly translate into an uncertainty in the above quantities. It is therefore necessary to study as carefully as possible all sources of uncertainties in the LEP energy calibration.

Flux-loop calibration

The basic relation between a charged particle's momentum transverse to the magnetic field in which it moves and the magnitude of this field is given by

$$p = \frac{R|q|}{B} \quad (2.1)$$

where q is the particle charge and R is the radius of the helix the particle will follow. The magnitude B of the magnetic field is then determined by monitoring the current going through calibrated dipole magnets.

The so-called 'flux-loop' calibration uses the voltage induced in coils ('flux loops') inserted in the lower poles of the dipole magnets, by the variation of the current going through the magnets. The flux loops of all dipoles in one curved section are connected in series. Prior to installation, the flux loops were calibrated using rotating coils with known properties placed on the theoretical central orbit of the magnets. A comparison of the voltage induced in the flux loops during a magnetic cycle from +2900 A to -2900 A with the voltage induced in the rotating coils then yields the effective surface of the flux loops. Using these effective surfaces one can relate the voltages induced in the flux loops to changes in the magnetic field, which (after comparison with the dipole currents) gives a calibration of the magnetic field generated by the dipoles. The precision of this calibration of the LEP beam energy is

around 20 MeV. In addition, the rotating coils provided information on the remanent magnetic field of the dipoles.

One stand-alone dipole is powered in series with the dipole magnets in the ring and equipped with a flip coil. The beam energy, derived from the flux measured by this coil, is continuously available. This energy value is used as a reference and is called the ‘field display’ value E_{FD} .

The flux-loop calibration does not take into account all effects on the magnetic field:

- The original calibration was performed without the vacuum pipe. A $7\text{ }\mu\text{m}$ thick nickel layer is used to bond this pipe to its surrounding lead radiation shield. Unfortunately this layer changes the magnetic field inside the vacuum pipe. In order to evaluate the consequences of this effect, the flux-loop calibration was redone for a dipole with the vacuum pipe and the nickel layer in place; the correction to be applied to the nominal energy value is $(+5 \pm 5)\text{ MeV}$ [22].
- The reference magnet has a core consisting of solid iron. The cores of the series dipole magnets on the other hand are made out of steel laminations; the spaces between the laminations are filled with cement mortar. This mortar suffers from a slow dehydration, causing both the surface of the flux-loop and the magnetic field length traversed by the particle beams to decrease. For the 1991 data taking period the flux loop results are corrected by a factor of $(+3 \pm 1) \cdot 10^{-4}$.
- A similar effect is caused by temperature variations. From June 1991 onwards, temperature sensors were installed on a set of eight magnets. The magnet temperature was found to vary between 17°C and 23°C . A correlation between the magnet temperature and the field display value was observed. This causes a relative correction to E_{FD} of $(1.0 \pm 0.25) \cdot 10^{-4}\text{K}^{-1}$ at 45 GeV.
- The remanent field of the magnets and their magnetic properties change with time. This introduces a nonlinearity in the energy calibration, which in turn results in a correction to be applied to the center-of-mass energy E_{CM} measured by LEP. For the 1991 data taking period this correction is found to be

$$\Delta E_{\text{CM}} = \alpha(2E_{\text{FD}} - 93\text{ GeV}), \quad \alpha = (-2.0 \pm 1.5) \cdot 10^{-3}. \quad (2.2)$$

The uncertainty on α is obtained by comparing the results of the proton calibration at 20 GeV and those of the resonant depolarisation measurements at 46.5 GeV, discussed below. The same uncertainty has been assumed for the 1990 data and for the 1991 pre-polarisation data.

Proton calibration at 20 GeV

A problem intrinsic to the flux loop calibration method is that other magnetic fields

(such as the earth magnetic field and the dipolar components of the orbit correctors and quadrupoles [23]) are not taken into account. It cannot therefore be regarded as yielding an absolute calibration. A calibration intended to overcome this problem is the so-called ‘proton calibration’. This method relies on the fact that at injection energy, protons are not yet ultra-relativistic. For a given particle momentum, the radius of the orbit described in a magnetic field is independent of the particle mass. The beam energy is measured indirectly by determining the velocity of protons circulating on the *same orbit* as the positrons.

Practical use of these facts is made by first injecting positrons. For a given magnetic setting (at injection energy) the central orbit (i.e. the orbit through the center of the quadrupoles) is determined. Subsequently the same is done for protons. The harmonic numbers associated with the RF oscillation frequency, $h_e = 31324$ and $h_p = 31358$ for the positrons and protons, follow from the circumference of the LEP ring. They are precisely known since a change by one unit would correspond to a change of 0.8 m in the circumference of the orbits [22]; this exceeds by far the precision with which the circumference of the LEP ring is known. The RF frequencies which lead to central orbits, f_e^{RF} and f_p^{RF} for the positrons and protons respectively, are measured. Approximating the positron velocity by $\beta_e = 1$, the velocity β_p of the protons is then given by

$$\beta_p = \frac{h_e f_p^{\text{RF}}}{h_p f_e^{\text{RF}}}. \quad (2.3)$$

The determination of f_e^{RF} and f_p^{RF} depends on the details of the transverse betatron oscillations of the circulating beams and can be performed with a precision of approximately 20 Hz [22]. The proton calibration method yields absolute beam energy values with an accuracy of approximately 1 MeV at the injection energy of 20 GeV. In order to obtain a calibration at 45 GeV one has to use magnetic measurements.

Resonant depolarisation

A more precise calibration can be performed by means of the so-called ‘resonant depolarisation’. This method makes use of the relation between the ‘spin tune’ ν (the number of spin precessions in one revolution) and the energy of the circulating beams:

$$\nu = \frac{g_e - 2}{2} \frac{E_b}{m_e c^2} \quad (2.4)$$

where $(g_e - 2)/2$ is the electron anomalous moment. Under ideal circumstances, an average transverse polarisation (i.e. a polarisation in the direction of the bending magnetic field) builds up spontaneously through the Sokolov-Ternov effect [24]. In the case of LEP the maximally attainable degree of polarisation was theoretically predicted to be 92.4% [25].

Resonant depolarisation occurs if a depolarising source is in phase with the spin precession. For a depolarising frequency f_{dep} smaller than the revolution frequency

$f_{rev} = f_e^{\text{RF}}/h_e$ this corresponds to the condition

$$\nu - \text{Int}(\nu) = f_{\text{dep}}/f_{rev} \quad (2.5)$$

where $\text{Int}(\nu)$, the integer part of the spin tune, is known unambiguously due to the precision obtained in other calibration methods; it therefore does not affect the precision of the present method.

The degree of polarisation P of the beam is measured by scattering a circularly polarised beam of laser light off the electron beam. The angular distribution of the Compton scattered photons depends on the electron helicity. The photons are detected 247 m from their interaction region in a tungsten calorimeter with a plane of horizontally running silicon strips of 2 mm pitch. Per ‘data point’ approximately $2.5 \cdot 10^6$ photons are used. The shift of the average y position of the backscattered photons upon a reversal of the laser beam helicity from $-\xi$ to ξ is a measure of the beam polarisation:

$$\Delta\langle y \rangle = \kappa \xi P, \quad \kappa = (500 \pm 30) \mu\text{m}, \quad (2.6)$$

where $|\xi|$, a quantity intrinsic to the laser apparatus, has been found to be limited to ~ 0.85 [25]. The laser scattering method enables a measurement of the degree of polarisation with a systematic uncertainty of $\Delta P = (0.5 + 12P)\%$.

A depolarising source is then introduced by letting a ‘kicker’ magnet produce an oscillating magnetic field in the horizontal plane. By slowly varying the oscillation frequency, a continuous monitoring of the degree of polarisation yields the precise frequency at which resonant depolarisation occurs. Using this method one finds a correction of

$$E_{\text{CM}} - 2E_{\text{FD}} = (61 \pm 5.3) \text{ MeV} \quad (2.7)$$

at a nominal energy of 93 GeV for the 1991 data taking period [22]. The main contribution to the uncertainty in this number (3.7 MeV) is due to the observed spread in the polarisation measurements.

The first polarisation measurement was performed in the middle of August 1991; strictly speaking the correction (2.7) is only valid for the physics data taken after this time. Fortunately, almost all of the off-peak data (which are most significant in the determination of the Z mass and its total decay width) were taken in the period following this first absolute calibration. In addition, polarisation was achieved at one energy point only. The aforementioned calibration techniques were used to obtain the corrections at the other energy points.

Other effects

Due to tidal effects, the LEP ring periodically changes shape. This affects the central LEP orbit. The length of the orbit used in the actual data taking is however fixed by the fact that the RF frequency f_e^{RF} remains constant. This means that the path followed by the particle beams through the magnets keeps changing. For an expansion of the ring there is an effective increase of the bending strength, implying

a decrease of the equilibrium LEP beam energy. The effect was measured using the resonant depolarisation method in a dedicated experiment, and was found to agree well with theoretical predictions. As a result, time dependent corrections of up to 18.5 MeV are applied [26].

The above-mentioned calibrations apply to the average beam energy in the arcs. An additional complication is caused by the radiation losses of the circulating beams. The RF cavities used to replenish these radiation losses are placed on both sides of IP2(L₃) and IP6(OPAL). The fact that two RF frequencies are used to reduce the power consumption by the cavities leads to an average phase error (for the higher frequency) of -21.4° for both the e^+ and e^- beams. This in turn leads to a difference of $\Delta E_{\text{CM}} = 16.1$ MeV for IP2. The effect is somewhat reduced by the fact that part of the RF cavities are superconducting and need only the lower RF frequency. The net correction for IP2 is 12.7 MeV. The uncertainty due to RF effects (tripped cavities, phase errors) is estimated to be 1.5 MeV. For an extensive account of all known uncertainties see reference 22.

Calibration results

For every beam energy setting the center-of-mass energy is affected by four sources of error:

- an absolute energy uncertainty ΔE^{abs} : this is dominated by the most precise calibration measurement and its extrapolation to the setting energy under consideration;
- a nonlinearity error ΔE^{nl} as defined in equation 2.2;
- a setting error ΔE^{set} : a random point-to-point energy error [27];
- a nonreproducibility error ΔE^{rep} : even with the same nominal machine energy settings the reproducibility of the beam energy is not perfect. This is partly due to temperature variations and RF instabilities; another part is presumably intrinsic and is estimated from the spread observed in the resonant depolarisation measurements. This contribution to the energy uncertainty decreases as the number of fills (and thus the number of energy measurements) increases.

The total relative energy uncertainty can be written as

$$\frac{\Delta E}{E} = \Delta E^{\text{abs}} \oplus \frac{|E^{\text{ref}} - E|}{E} \Delta E^{\text{nl}} \oplus \Delta E^{\text{set}} \oplus \frac{1}{\sqrt{\bar{n}_{\text{fill}}}} \Delta E^{\text{rep}} \quad (2.8)$$

where $\oplus$ denotes a summation in quadrature of the various sources; $\bar{n}_{\text{fill}}$ is the (luminosity weighted) number of fills taken at a given energy point. Table 2.1 summarises the uncertainties for the physics data taken in 1990 and 1991. It should be noted that for the first two errors there exists full correlation between the different energy points; for the uncertainties due to the last two sources there is no correlation between different energy values.

	1990	1991 (before Aug. 14)	1991 (after Aug. 14)
$\Delta E^{\text{abs}}(\times 10^{-5})$	29.0	20.0	5.7
$\Delta E^{\text{nl}}(\times 10^{-3})$	1.5	1.5	1.5
$\Delta E^{\text{set}}(\times 10^{-5})$	3.0	3.0	3.0
$\Delta E^{\text{rep}}(\times 10^{-4})$	1.0	1.0	1.0

Table 2.1: Contributions to the LEP energy uncertainty per energy point.

2.2 The L₃ Detector

The L₃ collaboration has optimised its detector for the best possible precision on electrons, muons and photons [13]. This goal has determined the overall design features of the L₃ detector (fig. 2.2). Going radially outward from the interaction point one encounters the following detector components:

- a small central tracking detector (CTD);
- a compact electromagnetic calorimeter (ECAL) based on the technique of total absorption in large arrays of Bismuth Germanate (BGO) crystals;
- a scintillation counter system providing precise timing signals;
- a hadronic calorimeter (HCAL) consisting of layers of depleted uranium acting as absorber, interspersed with wire chambers serving as active layers;
- a muon chamber (MUCH) system consisting of three layers of large and precise drift chambers.

In addition, the detector contains:

- a small angle luminosity monitor; like the ECAL it uses BGO crystals for performing electromagnetic energy measurements;
- a triggering and data acquisition system based primarily on FASTBUS techniques.

In order to enable precision momentum measurements in the muon chamber and in the central tracking detector, the whole detector is surrounded by a 7800-t octagonally shaped solenoid magnet providing a uniform axial magnetic field of approximately 0.51 T along the beam line. Figure 2.2 presents a general overview of the L₃ detector and its components.

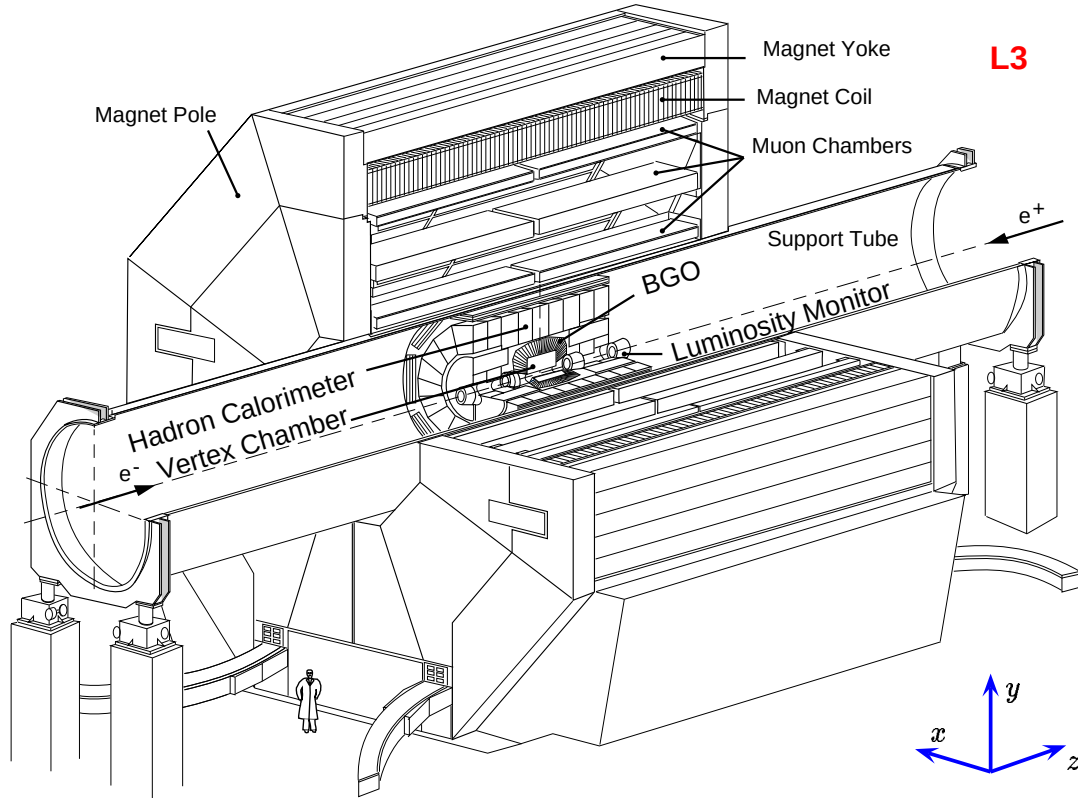


Figure 2.2: *The L_3 detector.*

Below follows a description of each of the above detector components. The (right-handed) L_3 coordinate system that is used in the rest of this thesis, has its origin in the center of the central tracking detector; its z axis lies along the direction of the e^- beam, its y axis lies in the positive vertical plane so that the x axis points towards the center of the LEP ring.

2.2.1 The central tracking detector

The design goals of the central tracking detector (CTD) are the determination of the origin, momentum and charge sign of charged particles emerging from the interaction point, and a reconstruction of secondary vertices for particles with lifetimes exceeding 10^{-13} s [13]. This subdetector consists of the following parts:

The time expansion chamber (TEC)

This chamber has an inner radius of 9 cm formed by a beryllium tube and an outer radius of 46 cm formed by an aluminium cylinder. Its sensitive length is 98 cm.

It is divided radially into an inner and an outer chamber; in the ϕ direction the inner chamber is divided into twelve, the outer chamber into 24 sectors. Electrons originating from ionisations caused by passing charged particles are drifted in a homogeneous electric field of 0.9 kV/cm towards a plane of anode wires, parallel to the beam direction and kept at a nominal voltage of 2.9 kV. Grouping the electron signals above a certain threshold then gives for each anode wire so-called ‘hits’. Each ϕ sector contains either 8 (inner chamber) or 54 (outer chamber) anode wires. A ϕ region corresponding to one inner sector and two outer sectors is shown in figure 2.3.

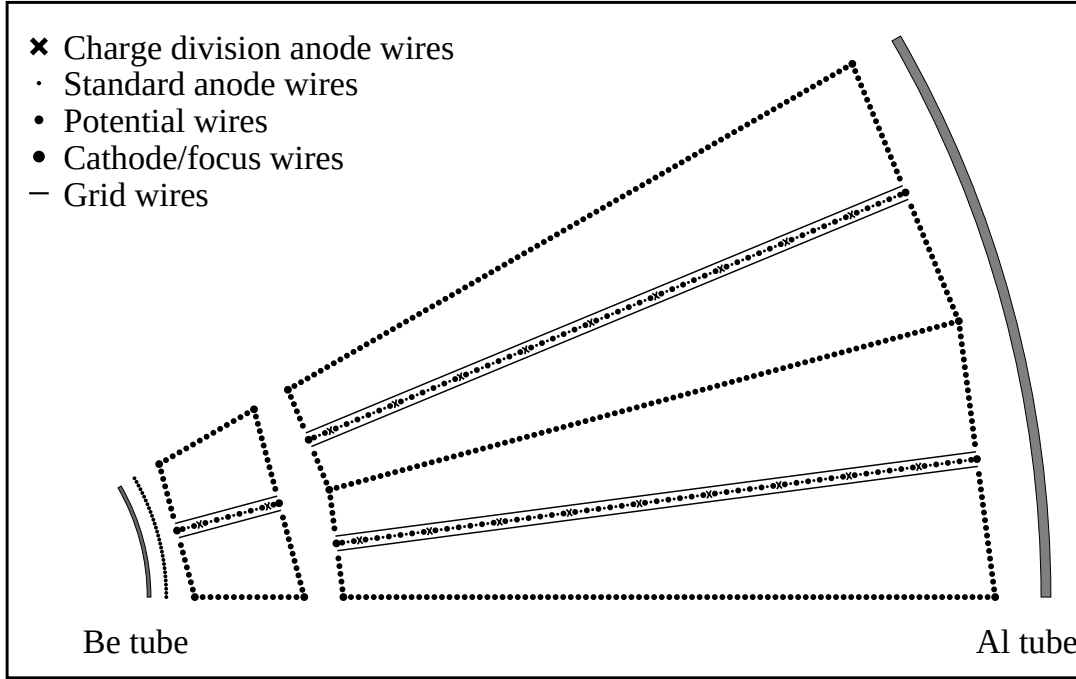


Figure 2.3: A *TEC* ϕ region viewed in the $x - y$ plane.

Close to the anode plane there is a plane of grid wires which are held at ground level. These ensure both the homogeneity of the field in the drift region and a high gas amplification in the region between the grid plane and the anode plane (the detection gap). The anode signals are read out on one end of the detector by 100 MHz Flash Analog to Digital Converters (FADC).

After conversion of the measured drift time into a drift distance, there is still a left-right (LR) ambiguity w.r.t. the anode plane. This can usually be resolved by combining the drift distance information coming from all anode wires associated with the track. In addition, in each grid plane of each outer chamber sector 14

groups of 5 grid wires are read out to further help resolving these LR ambiguities.

The chambers are filled with a mixture of 80% CO₂ and 20% isobutane $i\text{C}_4\text{H}_{10}$, at a temperature of 291 K and a pressure of 1.2 bar. This gas mixture has a low diffusion coefficient; it enables the use of low electric fields, thus decreasing the drift velocities (about $6\mu\text{m}/\text{ns}$) and improving the position precision. A single-wire precision of approximately $60\mu\text{m}$ is obtained [28].

In each ϕ sector the pulse height information for eleven of the anode wires (two in the inner chamber, nine in the outer chamber) is read out at both sides of the detector. A comparison of the pulse heights on both wire ends, called charge division (CD), then yields information on the z coordinates of the track created by the passing particle. The single-wire precision in the z direction is of the order of 6 cm [28].

The Z chamber

Since the precision that can be obtained using the CD information only is rather poor, additional measurements are needed in order to improve the z determination of charged tracks.

The Z detector serves this purpose. It consists of two cylindrical proportional chambers each with two layers of cathode strip readout having a pitch of 4.4 mm. Both chambers contain one layer of z -measuring strips (called z and z' layers) and one layer of strips installed at an angle of 70.1° with respect to the beam axis (called u and v layers). The u and v layers are crossed in order to provide additional stereo information. The Z chamber covers the polar angular range $45^\circ < \theta < 135^\circ$. For $\mu^+\mu^-$ final states the position accuracy provided by the Z chamber was measured to be $320\mu\text{m}$ [29].

Calibration and monitoring

For drift distance measurements, a precise calibration of the drift velocity is of utmost importance. This is performed by Plastic Scintillating Fiber (PSF) ribbons mounted onto the outer cylinder of the TEC. Each sector contains one ribbon; each ribbon contains 143 fibers running along the z direction. The tracks reconstructed in the TEC are associated with signals from the fibers. Plotting (per anode) the average measured drift time versus the position measured by the fibers then yields the (average) drift velocity.

To first order, the dependence of the drift velocity v_d on temperature T , gas pressure p and electric field E is given by $v_d \sim ET/p$ [30]. There is, however, also a strong dependence on the gas decomposition and the purity of the gas; it is therefore important that these conditions are kept as stable as possible. The TEC gas system, operating in a closed loop, is designed to keep the changes in gas pressure below 1 mbar and the changes in gas decomposition below 0.07 %. The cooling system has to perform to better than 0.1 K and the design of the high voltage is such that variations in the electric field are less than 0.1%.

2.2.2 The electromagnetic calorimeter

The electromagnetic calorimeter (ECAL) consists of large arrays of BGO crystals which transform electromagnetic energy depositions into a scintillation light signal. BGO was chosen because it has a very short radiation length ($X_0 = 1.12$ cm) thus permitting the construction of a very compact detector. For an extensive account of the properties of BGO see ref. 18. For electromagnetic energy depositions, the energy precision (in percent) obtained in the barrel region of the detector can be parametrised by [31]:

$$\frac{\sigma(E)}{E} = \sqrt{\left(\frac{1.54}{\sqrt{E}} + 0.38\right)^2 + (1.18)^2 + \frac{0.25}{E^2}} \quad (E \text{ in GeV}). \quad (2.9)$$

For energies above 1 GeV the angular precision is 0.5° [18].

The barrel part of the ECAL covers the polar angular range $42^\circ < \theta < 138^\circ$. It consists of 7680 crystals shaped as truncated pyramids, all pointing to the interaction region. For test beam purposes the barrel was divided into two half-barrels; in each of these the crystals are arranged in a 160×24 matrix in the $\phi - \theta$ plane; in the ϕ direction there is a small depointing offset (10 mrad) in order to improve the hermeticity of the detector.

The length of the crystals is 24 cm ($\approx 22X_0$). The front faces have dimensions of $(2 \times 2)\text{cm}^2$; at the back the dimensions vary (along the θ direction) between (2.6×2.6) and $(2.9 \times 2.9)\text{cm}^2$. In order to prevent permanent mechanical pressure from destroying the crystal properties [32], each crystal is housed in its own cell in a carbon fibre structure and kept against the front end of this cell by a spring-loaded holding device.

Each crystal is read out by two 1.5 cm^2 photodiodes mounted on its back face. The photodiode output signals are first preamplified and passed through a shaping circuit. The peak voltage of the resulting signal is used (after subtraction of the so-called ‘pedestal’ value, the average noise level) as a measure of the energy deposited in the crystal. Generally, crystals are read out only if the energy deposit exceeds a certain threshold, normally set to 2 MeV; however, one out of each 60 crystals connected in a token ring is always read out. This so-called ‘unsparsed crystal’ enables the detection of periods during which, for some reason, a ring was not included in the readout.

During the shutdown between the 1990 and 1991 data taking periods, additional arrays of BGO crystals (so-called ‘endcaps’) were installed in the forward and backward regions. Each of these endcaps contains 1527 crystals, in a more complex structure than in the barrel. Their polar angular coverage is $11.6^\circ < \theta < 38^\circ$.

Calibration and monitoring

In order to reach the BGO design precision, careful calibration and monitoring is of utmost importance:

- The reference temperature for the energy calibration is 18°C. This temperature was chosen as low as possible (to give the largest output signal) but safely above the dew point where condensation might damage the readout electronics. The calibration factors at this reference temperature were determined in a test beam at CERN before the installation of the barrel detector. This calibration sets the absolute energy scale.
- Since the light yield of BGO (i.e. the amount of light output generated per unit of energy deposited) has a temperature gradient of -1.55%/K the temperature is required to be stable within a few tenths of a degree. Temperature sensors are installed on every twelfth crystal and provide a temperature measurement every thousandth event. Corrections for temperature variations are applied offline.
- The pedestal value for each readout channel is determined before every LEP fill.
- The light yield as measured by the photodiode depends on the position of the energy deposition along the major axis of the crystal. Since the longitudinal shape of the shower generated by an electromagnetically interacting particle depends on the particle energy, this variation introduces a non-linearity in the relation between the output voltage from the electronics and the energy actually deposited in the crystal.

The variation of the measured light yield, as a function of the distance of the energy deposition to the photodiode, was measured using cosmic muon rays which generate a minimum-ionising signal in the BGO. The dependence of the measured output on the position of the muon track along the length of the crystal was determined. On the basis of the observed light yield variation, an energy-dependent correction factor was obtained and applied to the measured output signal.

- Finally, long-term variations in the measured light yield (i.e. in the above-mentioned absolute energy scale) need to be taken into account, as well as short-term effects caused e.g. by radiation damage to the crystals. In order to monitor these variations, a daily ‘Xenon’-run determines the output from the photodiodes for a known input coming from xenon flash lamps and fed into the crystals with the help of optical fibers. This output is then compared with the output obtained in the test beam, thus measuring the variation in the light collection efficiency.

A detailed account of the various calibration procedures can be found in ref. 18.

2.2.3 The hadron calorimeter

In the L_3 detector, the energies and directions of hadronically interacting particles are measured by means of a hadron calorimeter placed just outside the electromagnetic calorimeter. It consists of a barrel calorimeter, a forward-backward system and a muon filter covering 99.5% of the full solid angle (see figure 2.4). It is a fine sampling calorimeter made out of absorber layers interspersed with layers of gas-filled proportional wire chambers (PWC). Due to the space constraints imposed by the muon chamber system the detector had to be built as compact as possible while retaining its stopping power for hadrons. A high stopping power ameliorates the energy precision of the calorimeter and reduces the hadron background in the muon chambers. To achieve the best absorption within the volume allowed by the muon chambers uranium was used as the absorber medium. The energy precision as determined in hadron test beams is $\sigma(E)/E = (55/\sqrt{E} + 5)\%$ (E in GeV) [33]; the angular precision for hadronic jets in physics data is 37 mrad as determined from a sample of two-jet events [34].

The barrel detector

The barrel system covers the polar angular range $35^\circ < \theta < 145^\circ$. It consists of nine rings each containing sixteen modules. The outer radius of the detector is 1975 mm. The middle three rings have an inner radius of 885 mm; the other rings have an inner radius of 979 mm. The total length is 4725 mm. All modules consist of a sandwich of 5 mm thick plates of uranium and layers of proportional wire chambers filled with a mixture of 80% Ar and 20% CO_2 . The chambers are read out by charge integrating ADCs with a sensitivity of 50 fC (~ 25 MeV) per count. A 54 mm thick stainless steel plate shields the BGO photodiodes from the uranium radioactivity. The amount of material encountered by particles originating from the interaction region varies between 6 and 7 nuclear absorption lengths.

The chambers operate at a high voltage of 1.6 kV giving a gas gain of about 10^4 . In successive chamber layers the wires are stretched alternately in the z direction and the ϕ direction in order to improve the spatial information. The long (short) modules contain 58 (53) layers of chambers. The readout of the (in total) 371,764 wires is grouped in so-called ‘towers’, reducing the number of readout channels to 23,040; the granularity decreases with the distance from the beam axis. In the ϕ direction each of the towers points to the interaction region; in the z direction the towers have a constant width.

The forward-backward calorimeter

The forward-backward system covers the angular ranges $5.5^\circ < \theta < 35^\circ$ and $145^\circ < \theta < 174.5^\circ$. It consists of three modules on both sides, each of which split vertically into half-modules for practical reasons. The half-modules consist of 77 (HC1), 27 (HC2) and 23 (HC3) chamber layers respectively. The wires in successive layers

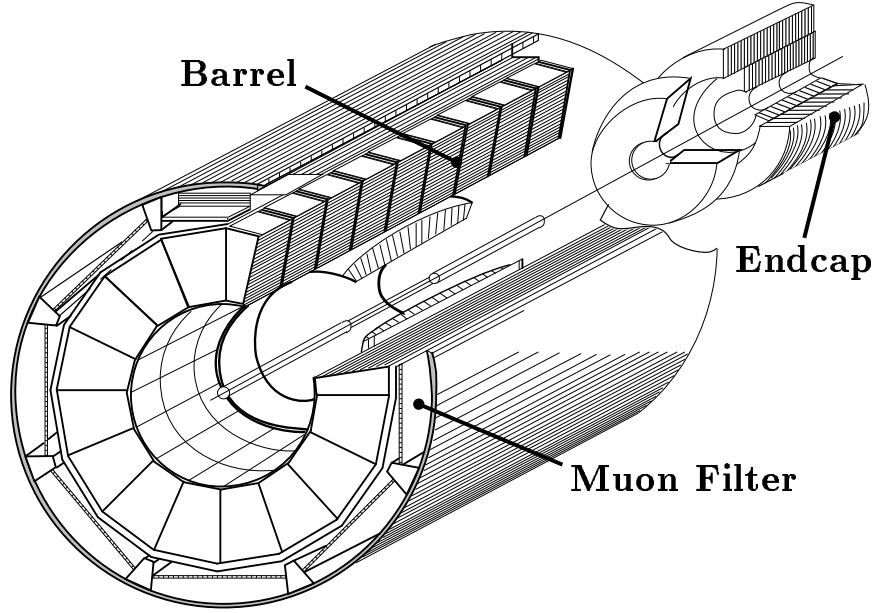


Figure 2.4: *Perspective view of the L_3 hadron calorimeter.*

(‘U’ and ‘V’ layers) are rotated by 22.5° with respect to each other. The system is subdivided into a fine sampling part (with 5 mm thick uranium plates) in front of a coarse sampling part (10 mm thick plates). The first three layers of plates in the fine sampling part are replaced by steel plates in order to shield the ECAL endcaps. The plates are arranged such that the gaps between them never point towards the beam axis; the gaps in successive layers are at different places so as to ensure a good hermeticity. The readout of the 54,140 wires is again organised in towers, such that all towers point towards the interaction region in the θ direction, with a granularity of $\sim 2^\circ$. The azimuthal granularity provided by the U and V layers is 22.5° .

The muon filter

The muon filter provides another absorption length of material in front of the muon chamber system. It consists of eight octants situated just inside the support tube. The sandwich structure in this case consists of five proportional chambers, operated with the same gas mixture as the barrel and endcap HCAL parts, interleaved with six 10 mm thick brass absorber plates. To maximise the stopping power for hadrons, absorber plates without chamber layers fill up the remaining space inside the support tube. Each octant has a length of 4 m and a width of 1.4 m. In total 624 chambers along the z direction, with a total of 8064 channels, are read out. The high voltage applied is somewhat higher than in the barrel and endcap systems (1.8 kV), such that the higher gas gain gives a larger signal for minimum ionising particles. Charge division is applied for the three middle layers of each octant.

Calibration and monitoring

The uranium absorber medium of the barrel and endcap systems provides a calibration mechanism through the photon spectrum generated by its natural radioactivity. The calibration method makes use of the fact that the resulting pulse height spectrum changes with the gas gain and the preamplifier behaviour. For this purpose, a high statistics ‘reference’ spectrum is recorded in a monthly calibration run, yielding gain accuracies of 0.3% per readout channel. In addition, during the startup or end period of each run a spectrum is recorded; since the pulse height varies linearly with the gas gain a comparison of the spectrum to the reference spectrum yields the instantaneous gain relative to the reference spectrum. The point-to-point precision one obtains in this way is better than 2% [35]. The gain variations are taken into account in the offline reconstruction.

The gas gain depends on gas temperature and decomposition, the high voltage applied and the amplifier gain. In order to obtain the best possible precision, these are kept as constant as possible.

2.2.4 The muon chamber system

The L_3 muon chamber system (MUCH) consists of large drift chambers, placed inside the L_3 magnet volume. The system consists of two ferris wheels each containing eight octants, placed at radial distances of 2530, 3985, and 5440 mm from the beam line respectively. The three layers of so-called P chambers cover a polar angular range $45^\circ < \theta < 135^\circ$ and provide a momentum precision of 2.5% for 45 GeV muon tracks. Tracks measured in only two layers have a precision of approximately 20%.

The z coordinate measurement of the track is performed by four layers of Z chambers, arranged on the top and bottom of the inner and outer chamber layers. Each of these layers in turn consists of two sublayers, staggered in order to resolve LR ambiguities. As the information provided by the muon chamber system, although taken into account, plays only a minor role in the reconstruction of hadronic final states, no detailed description of this detector is presented here. For an extensive account of the muon chamber system see for instance reference 36.

2.2.5 The luminosity monitor

The L_3 luminosity monitor is designed to precisely measure the LEP luminosity in the L_3 interaction point, by means of a measurement of the rate of small angle Bhabha events. The two parts of this subdetector are nominally located at $z = \pm 2.765$ m from the interaction point and cover a polar angular range $24.7\text{mrad} < \theta < 69.3\text{mrad}$. The actual geometry of the detector is measured several times per year by a CERN survey group. Each array consists of 304 BGO crystals, arranged in cylindrically symmetric layers and azimuthally subdivided into sixteen sectors. A schematic view of the luminosity monitor is given in figure 4.1. The crystals are

26 cm ($\sim 24X_0$) long; the dimensions of the front faces vary between $(1.5 \times 1.5)\text{cm}^2$ and $(1.5 \times 3.0)\text{cm}^2$. The crystals are read out by the same system (with minor modifications) that is used to read out the ECAL barrel and endcap crystals.

During physics runs but before the beams are collided, instabilities can result in a beam loss which, if occurring near the L₃ detector, could severely damage the luminosity monitor BGO crystals. Therefore each array is split vertically into two parts which can be moved apart by means of a hydraulic system; the positioning accuracy in the closed position is $10\text{ }\mu\text{m}$ [13].

A daily calibration run using the known input from LEDs placed opposite the photodiodes enables a monitoring of the short-term variations of the crystal light yield, such as those caused by radiation damage. The long-term light yield variations are taken into account by performing, several times per year, a calibration using the luminosity data collected during the preceding period. At 45 GeV, an energy precision of 2% is achieved [37].

2.2.6 The scintillation counters

The study of dimuon events is usually complicated by the fact that there are always cosmic muons present. If one such muon traverses the detector close to the interaction point it becomes impossible, on the basis of the event topology alone, to distinguish it from a genuine dimuon event originating at the interaction point. For this reason, a set of 30 plastic scintillation counters provide a timing signal that enables the rejection of most of the cosmic background.

The scintillation counters are mounted on the inner side of the HCAL barrel. They cover the polar angular range $34^\circ < \theta < 146^\circ$ and, since they follow the shape of the HCAL barrel modules, they are wider at the ends (182 mm) than in the center (167 mm); 93% of the full azimuthal range is covered. The radial distance from the interaction point to the 1 cm thick counters is 885 mm in the central region and 974 mm in the forward and backward regions.

The photomultipliers used to amplify the scintillator signals have a transit time jitter of 0.6 ns for 10 photoelectrons (to be compared with an average of 80 photoelectrons per minimum ionising particle) and are read out by TDCs with a time precision of 50 ps. Each scintillator is read out on both sides. The time difference between the signals from the ‘near’ and the ‘far’ counter yields the position (in the z direction) where the scintillator was hit; this serves to further improve the time precision. Thus a precision of approximately 0.6 ns is achieved [13].

2.2.7 Trigger and data acquisition

Out of every $n_{\text{bunch}} \cdot f_{\text{rev}}$ bunch crossings per second, one expects (with nominal currents of $500\text{ }\mu\text{A}$ /bunch and under normal LEP data taking conditions) a ‘physics’ event rate of the order of 1 Hz. In order to maintain a reasonable tape writing rate,

a complex trigger and data acquisition system was set up. This reduces both the detector ‘dead time’ (the time during which the detector cannot take data) and the subsequent amount of offline computing. Each of these systems consists of several logical levels.

The data from all subdetectors is digitised and sent to subdetector FASTBUS crates. Together with the data, each subdetector sends the current event number, as provided by the trigger control, in order to prevent event desynchronisation between the subdetectors. All data are stored onto IBM cartridge tapes via a central FASTBUS crate connected to a VAX 8800. Their subsequent reconstruction (see chapter 3) is handled by a central data production system running on an IBM main-frame system. The organisation of the data production is discussed in some detail in reference 38.

The trigger level-1 is designed to make a fast decision (i.e. before the next bunch crossing, $22(11) \mu\text{s}$ for $4 \times 4(8 \times 8)$ operation) whether or not to accept an event for further processing. If this decision is negative, the whole system is reset so that it is ready to process the next event; if it is positive, the subdetectors receive a signal which tells them not to accept any following event until the next reset is received. Only then does buffering and transfer of the data take place.

The level-1 decision is itself made up of the logical OR of four independent sub-triggers: the energy trigger, the muon trigger, the TEC trigger and the scintillator trigger. Below follows, in some detail, a description of each of these:

Energy trigger

This trigger uses the signals from all calorimeters (i.e. ECAL, HCAL and luminosity monitor). For the ECAL, every 30 crystals give one trigger signal; for the HCAL every module, except the outermost ones, gives two signals corresponding to the first interaction length and the rest of the calorimeter respectively. For the luminosity monitor every ϕ sector gives one signal. In order to reduce the amount of trigger data, all ECAL and HCAL signals corresponding to the same θ or ϕ value are added. Based on the resulting signals the following triggers are defined:

- a cluster trigger looking for energy depositions in different detector layers in the same direction;
- a total energy trigger requiring a minimum total energy or a minimum energy in the barrel region of the detector;
- a single-photon trigger asking for a single isolated energy cluster in the barrel of the ECAL;
- a hit counting trigger requiring a minimum number of calorimeter hits;
- a luminosity trigger requiring a back-to-back correlation between the two opposing luminosity subdetectors.

As the calorimeters are sensitive to noise from various sources, a frequent calibration of the trigger analog sum channels is necessary; such a calibration is normally performed after a calibration of the readout channels.

Muon trigger

The muon trigger uses all muon P chamber and Z chamber signals. For the P chambers, all wires corresponding to a chamber cell are grouped together; for the Z chambers, adjacent wires in each double plane constitute one ‘trigger cell’. The information from all trigger cells is encoded and matched against predefined ‘roads’ (defining all possible tracks with $p_T > 1$ GeV). The matches in all planes (both P and Z) are stored. Using the coincidence of matches in different planes the following triggers are set up:

- a single muon trigger, requiring at least one coincidence of all P layers and one coincidence of all Z layers;
- a dimuon trigger, asking for two coincidences, each involving at least two P layers and two Z layers; in addition, these two coincidences are required to be sufficiently acoplanar and in different octants;
- a small angle muon trigger, requiring two coincidences of the MI P and Z chamber layers, in opposite halves of the detector.

TEC trigger

The TEC trigger uses all of the 14 outer chamber wires which carry LR information. Using these wires to solve the LR ambiguities and dividing every half-chamber in two – according to the drift time – one obtains 96 ϕ bins in total. On the basis of the resulting 14×96 points a track search is performed, comparing hit patterns with predefined roads (as in the case of the muon trigger) and taking into account possible inefficiencies and spurious hits. Three categories of tracks are looked for:

- tracks with polar angles $37^\circ < \theta < 143^\circ$ and transverse momenta $p_T > 600$ MeV; these tracks are normally confined within one and the same ϕ bin;
- tracks with $37^\circ < \theta < 143^\circ$ and $p_T > 150$ MeV; these tracks can occupy up to three adjacent ϕ bins;
- tracks with $25^\circ < \theta < 155^\circ$ and $p_T > 100$ MeV; such tracks are identified with the help of the seven innermost wires only.

Given the long drift times occurring in the outermost wires, the forward tracks are searched for first; this is followed by the search for central tracks. The final trigger decision depends on the total number of tracks, the number of clusters of contiguous tracks, and the number of pairs of nearly coplanar tracks.

Scintillator trigger

The signals from the 30 scintillation counters are used for this trigger. Its subtriggers are [39]:

- a scintillator multiplicity trigger: for this trigger the counters are grouped in pairs. A coincidence of at least five pairs (out of 16) is required, with at least two pairs separated by more than five pairs;
- a coincidence pattern trigger: this trigger uses the signals grouped in four quadrants. Depending on the background rates, either just a hit or a back-to-back coincidence is required. This trigger is always used in a logical AND with the muon trigger.

Upon a positive level-1 decision, the results of the level-1 calculations are transferred to the level-2 event buffers together with the TEC charge division information. Each trigger sends its trigger data together with the event number; this enables the level-2 to check the event synchronisation. The level-2 trigger takes on average ~ 8 ms/event and further reduces the event rate by a factor of two. It does so by examining more closely the energy distribution within the event and by using the three-dimensional TEC track information provided by the charge division wires.

The trigger level-3 constitutes the last online stage of event rate reduction. In the case of a positive level-2 decision, the level-3 avails itself of the level-1 and level-2 results, as well as of all the digitised data. The latter can only be used at level-3 because the Data Reduction Processors (DRP) used to reduce the amount of data coming from the TEC take 20 ms to fulfill their task. Since the level-3 trigger has much more time available than the lower level triggers it is able to perform more complex calculations. It is this fact which enables the level-3 processors to achieve an event rate reduction by another factor of two. Using the above staggered trigger level scheme, the level-1 trigger rate is only about 10 Hz under normal circumstances; the dead time is then approximately 5%. This dead time (about 6 ms) is mainly due to the long digitisation time in the ECAL readout.

2.2.8 Event simulation

Prior to their installation, various parts of the L_3 detector were tested extensively in test beams in order to study their energy and position precisions. However, the monochromatic test beams represent an environment which is often quite different from that encountered in genuine data taking conditions. In the case of hadronic events, for instance, the response of the detector to hadron jets is a property which cannot be determined from test beam data alone. A test beam does not provide the complexity of real events.

The high precision measurements aimed at in the LEP experiments require a thorough understanding of the detector response to various signals. In addition, the

complications due to fluctuations both in the physics signals (i.e. in the particle final states) and in the response of a complex detector cannot be dealt with analytically. For this reason one turns to Monte Carlo (MC) techniques [40]. The Monte Carlo method simulates ‘events’ which can be compared with experimental data.

The simulation of events proceeds in two steps. The first step makes use of so-called *event generators*. The method consists of generating particle configurations allowed by a physical model; these events are assigned a weight equal to the differential cross section for the model in question. Using these ‘weighted’ events one can calculate – again in a model dependent context – the cross section for a process based on any given kinematic and fiducial volume cuts. The importance of event generator programs lies in the fact that any physical model can be used as input; any known or hypothetical final state can therefore be studied. It is also possible to apply an ‘event rejection’ procedure such that the number of MC events kept in a particular region of a multiparticle phase space is proportional to the differential cross section in that phase space region. The ‘unweighted events’ thus obtained are generally used in detector studies.

The second step in the production of MC events uses the *detector simulation* programs. The input to these programs is the particle final state of a sample of unweighted events. The purpose is to simulate the response of the various subdetectors to the generated particle configuration. In this stage no event rejection is applied anymore. In order to obtain a correct simulation of experimental data it is necessary to properly take into account, for each particle separately, its lifetime and its interaction cross sections in the various detector materials. Where applicable, additional effects caused by the detector readout (notably Birks’ effect in the case of the ECAL and the HCAL) are taken into account as well.

For detector simulation purposes, all four LEP experiments use the GEANT3 [41] package. This program allows a detailed geometrical description of the detector; it is interfaced to subprograms describing interactions which are relevant for present-day detectors.

A proper treatment of all possible physics processes is most important in the case of calorimetric measurements. In tracking detectors, pattern recognition algorithms have to be applied. These algorithms normally allow for some inefficiencies, thus providing a relatively high degree of robustness. In the case of calorimetry, again some sort of pattern recognition has to be carried out; however, now each individual energy deposit can have a much larger weight. This is mainly due to the fact that the pattern is often only two-dimensional and fluctuations can be much larger. The above holds in particular for the calorimeters used in the L₃ experiment.

In order to provide a good representation of physics data, the various interactions between incoming particles and the detector materials must be described in great detail. The electromagnetic interactions are well understood; therefore in this case the only important limitations are those of computing time – generally this implies a cutoff in the energies of particles involved in a shower cascade below which inter-

actions are treated only collectively. The case of hadronic interactions is somewhat more complicated. This is because there is a very broad range of strong interaction types; many of the cross sections are poorly known, in which case one has to revert to parametrisations of experimental data[42]. The parameters involved are tuned using the results of test beam measurements [13]. A check on the validity of the results of the simulations is provided by the existence of several independent packages describing the same interactions. A good agreement is observed in most cases.

Chapter 3

Event Reconstruction and Selection

In order to determine the cross section for the process $e^+e^- \rightarrow q\bar{q}$ it is first necessary to identify the hadronic final states among the events recorded by the detector. Due to the finite precision and limited efficiency of the various subdetectors, and because other channels can look like hadronic events, this identification is never perfect. Therefore an event selection procedure is carried out, the aim of which is to correctly determine the *number* of hadronic events in a given sample. This is done by subjecting the event sample to certain selection criteria or cuts, based on the – supposedly – known shapes of both signal and background events. Such cuts are optimised to reject the background while retaining as much as possible of the signal.

The output signals from the various subdetectors form the ‘raw’ data. In order to identify the physical content of an event it is necessary to reconstruct the particles that originated from the e^+e^- interaction. This is done in two stages: first all subdetector signals are grouped into higher level objects; subsequently these subdetector objects are combined so as to yield global quantities referring to the particles in the event. In the following each of these reconstruction stages is considered in some detail.

3.1 Event Reconstruction

3.1.1 CTD track reconstruction

The aim of the reconstruction in the central tracking detector is to determine the momenta and charge signs of the charged particles in an event. The reconstruction of charged tracks in the CTD is a nontrivial matter due to the fact that the elements involved are three-dimensional; this leads to considerable pattern recognition problems.

The reconstruction in the CTD is done in two steps. In the first step, tracks

are formed from hit patterns and the track parameters in the transverse plane are determined. The second step aims at obtaining the polar angular information.

The search for tracks in the transverse plane starts from pairs of hits with similar drift distances to adjacent wires. Subsequently the track radius of curvature is calculated. This is done either with the help of the nominal interaction point as a hypothetical track point, or, if a hit belongs to two pairs, using only the hits themselves. Additional hits are then associated with this ‘tree’ on the basis of the distance to the track thus predicted. Finally, the track segments are split or combined, until an acceptable sample of tracks is found, i.e. until satisfactory circle fits to the track hits in the transverse plane are obtained. The LR ambiguities are resolved (as mentioned in section 2.2.1) by combining the segments reconstructed in the inner and outer chambers. A global track fit is then performed. Given the homogeneity of the magnetic field inside the small volume of the TEC a relatively simple circle fit is adequate. Figure 3.1 shows the distribution of the number of reconstructed tracks for the 1991 data and the corresponding Monte Carlo predictions¹. The discrepancies are caused by inactive TEC sectors in the real data. However they do not intervene in the event selection.

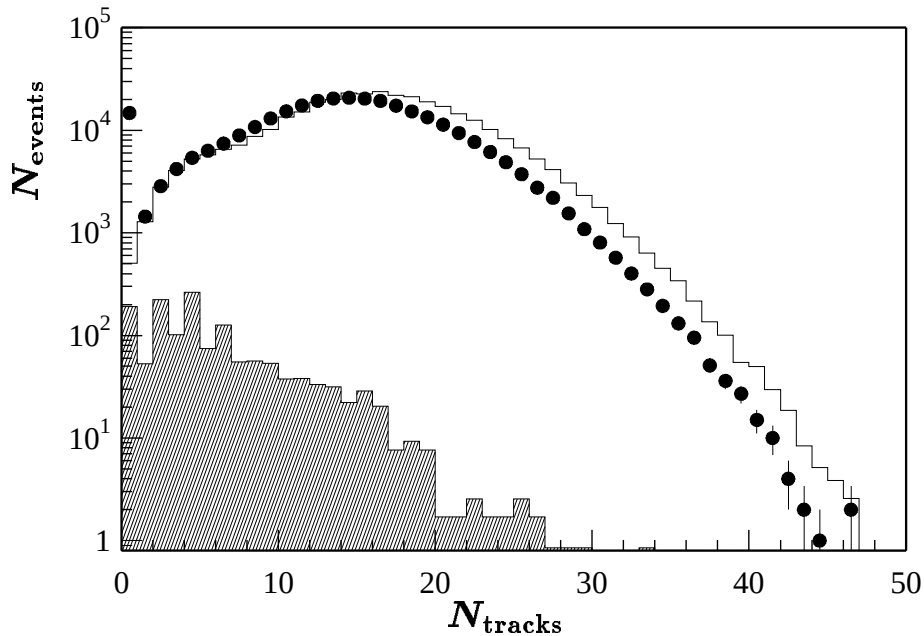


Figure 3.1: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the distribution of the number of reconstructed CTD tracks. The hatched area denotes the background predicted by the Monte Carlo.

¹The normalisation of the Monte Carlo histograms in this section is based on the total area covered by the plots.

The drift velocity v_d comes into play when converting from drift times to drift distances. A good transverse momentum (p_T) precision requires that changes in the instantaneous drift velocity be taken into account. These changes are the result of variations of the temperature, pressure and composition of the drift gas, and of variations in the electric field. Use is made of a database, which stores all time dependent information relevant for the operation of the CTD. The known dependence of v_d on the gas parameters allows the application of offline corrections.

The polar angle information is obtained by considering, for each track, the signals from the associated CD wires. This does not require pattern recognition. The pulse height information used in this part of the reconstruction is sensitive to noise and to effects from overlapping hits. As a result, the single wire precision in the z direction is orders of magnitudes worse than in the transverse plane: approximately 6 cm for low multiplicity events and 15 cm in the case of hadronic events. Additional z measurements are obtained from the Z chamber layers. Only signals which can be associated unambiguously with the track are considered; the track information already available is used to constrain the possible matches. For low multiplicity events only a ϕ window is used. Using the stereo information provided by combining the signals from different Z chamber layers, associations unambiguous in the ϕ direction are searched for. For high multiplicity events usage of the ϕ constraint alone becomes insufficient due to combinatorial problems; in this case also a window in the z direction (obtained by a fit to the CD signals) is used.

3.1.2 ECAL reconstruction

The purpose of the reconstruction of objects in the ECAL is to determine the energies and directions of particles interacting with the BGO material, and to classify them as either particles causing electromagnetic showers (electrons, positrons and photons), particles inducing minimum ionising signals (mostly muons) or particles with mainly hadronic interactions.

As a first step, the signal measured in each crystal is converted into an energy deposition. This involves a calibration which differs from crystal to crystal. To keep track of time dependent and crystal dependent variations (primarily the pedestal and the temperature variations), again use is made of a database. In the Monte Carlo events, one masks – per run – the crystals which were found to be inactive in the real data. As an illustration, figure 3.2 shows the total energy distribution measured in the ECAL for the 1991 data. A good agreement between the data and the corresponding Monte Carlo predictions is obtained.

The next step consists of grouping the energy depositions in the different crystals into so-called ‘clusters’; clusters are defined as contiguous regions of crystals with at least 10 MeV of energy deposit per crystal. Subsequently, the most important step is the search for local maxima (‘bumps’) in the energy deposit patterns of these clusters. A bump is formed either by a crystal having an energy deposit of at least

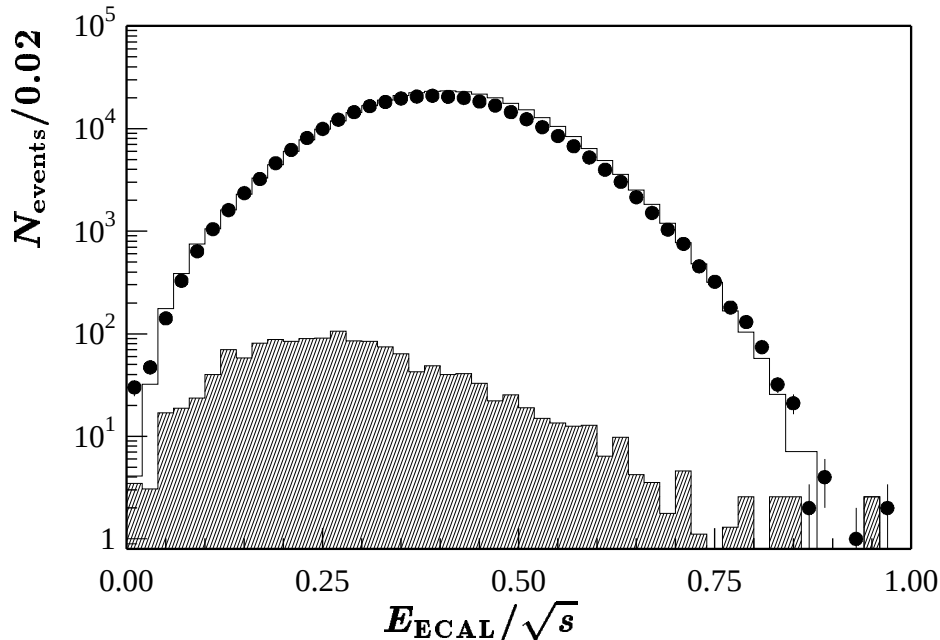


Figure 3.2: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the distribution of the energy measured in the ECAL. The hatched area denotes the background predicted by the Monte Carlo.

40 MeV or by a group of crystals containing at least one such crystal. In the latter case the most energetic crystal is called the ‘bump crystal’. If a crystal is equidistant to more than one bump crystal, it is assigned to the most energetic one. If bumps are overlapping, no attempt is made to determine the actual sharing of the energy between the bumps involved.

The best precision using the ECAL is obtained for electrons and photons. The main reason is that the resulting showers are usually well contained (the average *longitudinal leakage* being always smaller than 1% [18]). Hadronic interactions in the BGO lead to much larger fluctuations, because they are less well contained (the ECAL represents only one nuclear interaction length) and more complex. The approach taken therefore is to attempt precise energy and position measurements for electrons and photons only.

Bumps are classified as being of electromagnetic origin or not on the basis of the observed energy deposition pattern in the crystals belonging to the bump. The number of crystals taken into account depends on the bump energy. A χ^2 test using test beam information is applied to make the classification. The raw bump energy (the simple sum of the energies of all crystals associated with the bump) and the center of gravity of the bump are computed. Corrections to these quantities are applied *a posteriori*, based on the assumption of an electromagnetic energy deposition. The energy precision achieved for electromagnetic interactions is given

in eq. 2.9: this implies a precision better than 1.9% for energies above 2 GeV. The angular precision is a few mrad. The same corrections are applied in the case of non-electromagnetic bumps. In principle this is not a correct treatment; however, since the precisions in this case are typically an order of magnitude worse than for electromagnetic energy depositions, this has no practical consequences.

The ECAL readout is sensitive to electronic noise. This noise consists of two parts: an uncorrelated part due to random fluctuations in the readout electronics, and a correlated part caused by external signals. In general, the uncorrelated part does not cause too many problems and can to a large extent be avoided by considering only bumps consisting of at least two crystals. Correlated noise on the other hand can involve a very large number of channels; special algorithms are needed to reject this type of noise. This problem is relevant mainly for the 1991 data taking period, after a servo device was installed used for maintaining the alignment of the L₃ support tube with respect to the LEP beam. The ECAL readout electronics turned out to be quite sensitive to the electromagnetic field generated by the motor effectuating these movements. The noise problem is discussed more thoroughly in reference 18.

3.1.3 HCAL reconstruction

The pattern recognition algorithm applied in the HCAL is similar to, but more involved than the one used in the case of the ECAL. It starts from the individual tower signals. A tower or ‘hit’ is only taken into account if its energy deposit exceeds 9 MeV. This cut rejects most hits due to uranium noise: the rms of the noise for a single tower is only 2 MeV. This leaves on average 4 hits per barrel module due to the background from uranium noise [34]. In the conversion from the electronics signal to energy deposits, the time dependent effects (in this case pedestal and gain variations) are again inferred from database information. Figure 3.3 shows the HCAL energy distribution for the 1991 data, and the corresponding Monte Carlo prediction. One observes a small systematic discrepancy which is however corrected for later on by the jet energy calibration described in section 3.1.5.

Also the HCAL hits are then grouped into clusters. Here, clusters can consist of hits which are neighbours either in the radial or in the angular direction; some allowance is made for inefficiencies. Since the HCAL segmentation is three-dimensional, a more complex pattern recognition algorithm must be applied than for the ECAL. In principle however, the algorithm is again based on the association of nearby hits.

The pattern recognition also distinguishes hadrons from minimum ionising muons using the fact that the energy depositions caused by a traversing muon are very localised near its track.

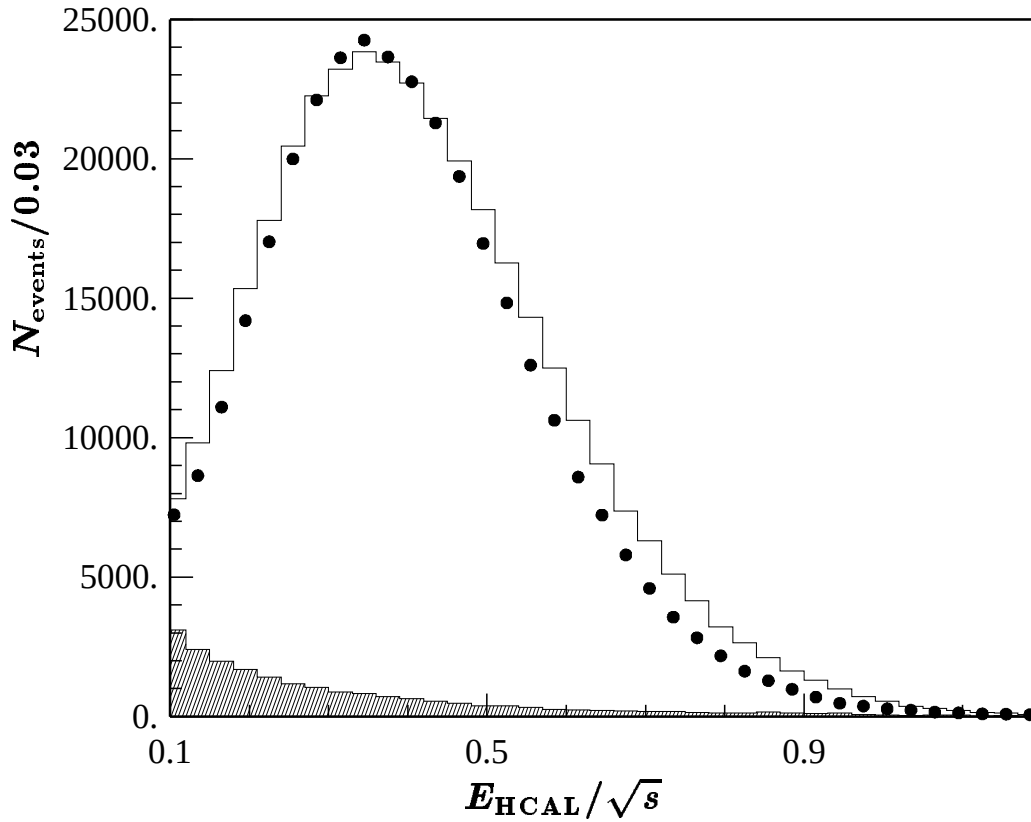


Figure 3.3: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the distribution of the energy measured in the HCAL. The hatched area denotes the background predicted by the Monte Carlo.

3.1.4 Muon chamber reconstruction

Given the good stopping power provided by the HCAL and the fact that the number of muon tracks in an event rarely exceeds two, the occupancy in the MUCH is generally low. This greatly facilitates the pattern recognition.

The search for tracks in the MUCH starts from the separate hit signals in the P and Z chambers. The first step is to combine the hit patterns in both of these chambers into segments.

Subsequently, each reconstructed Z ‘track’ is matched against all possible P segments in the same octant. Combinations of up to three P segments are considered. The final pattern recognition step is an attempt to identify tracks which cross octant boundaries. The best momentum precision is only obtained through a proper treatment of all systematic effects contributing to the hit coordinate measurements. Finally, in order to obtain an estimate of the momentum of the muon near the interaction point, corrections are applied for its average energy loss and for its multiple

scattering in the material traversed before entering the muon chambers. Details of the MUCH reconstruction algorithm can be found in reference 38.

3.1.5 Global reconstruction

The main aim of the event reconstruction is a detailed knowledge of the particle content of a given event. To reach this goal, it is often not sufficient to consider individually the information provided by the different subdetectors: it is necessary to combine all the available information.

There are several global reconstruction algorithms. They are all biased towards the optimisation of a particular result; which algorithm is most suitable depends on the specific analysis involved.

The algorithm used in this thesis uses information from the CTD, the ECAL, the HCAL and the MUCH, but depends mainly on calorimetric quantities. The first step is to associate the CTD tracks with the energy depositions in either the ECAL or the HCAL, on the basis of their angular proximity. Subsequently, bumps reconstructed in the ECAL are associated with clusters found in the HCAL. In this step, some use is already made of threedimensional information: ECAL bumps are not associated with HCAL clusters which have only energy deposited in the rear part of the HCAL; such energy depositions are unlikely to originate from the same particle. As a last step, muon tracks are associated with the closest energy depositions.

The resulting objects are the so-called ‘Smallest Resolvable Clusters’ (SRCs). On the basis of the subdetector information, each cluster is classified as being of electromagnetic, hadronic or of mixed origin. This classification is later used to improve the energy and position precision.

The clusters thus reconstructed are used to derive global event information. For hadronic events, the relevant quantities are the so-called particle ‘jets’, the physical origin of which is discussed in some detail in the next section. These jets are reconstructed from an energy ordered list of all clusters. Starting from the most energetic cluster, all clusters within a 30° cone are grouped into a supercluster called a jet. A jet axis is defined as the center of gravity direction of all the clusters assigned to the jet. Using this axis a new cone is defined in which one searches for additional clusters; a new jet axis is then calculated. This procedure is iterated until no more additional clusters are found. The left-over clusters are classified as being either single electrons or single photons, single hadrons or single muons. If such a cluster is separated by less than 20° from the nearest cluster in the nearest jet it is associated with this jet. The jets reconstructed in this way are required to have at least 7 GeV of energy. The jets are subsequently classified as being electromagnetic, single hadron or multihadronic. Figure 3.4 shows (for the 1991 data) the distribution of the number of jets thus reconstructed, and the corresponding Monte Carlo prediction; a satisfactory agreement is observed.

The energy assigned to each jet is based on calibrations performed for each of

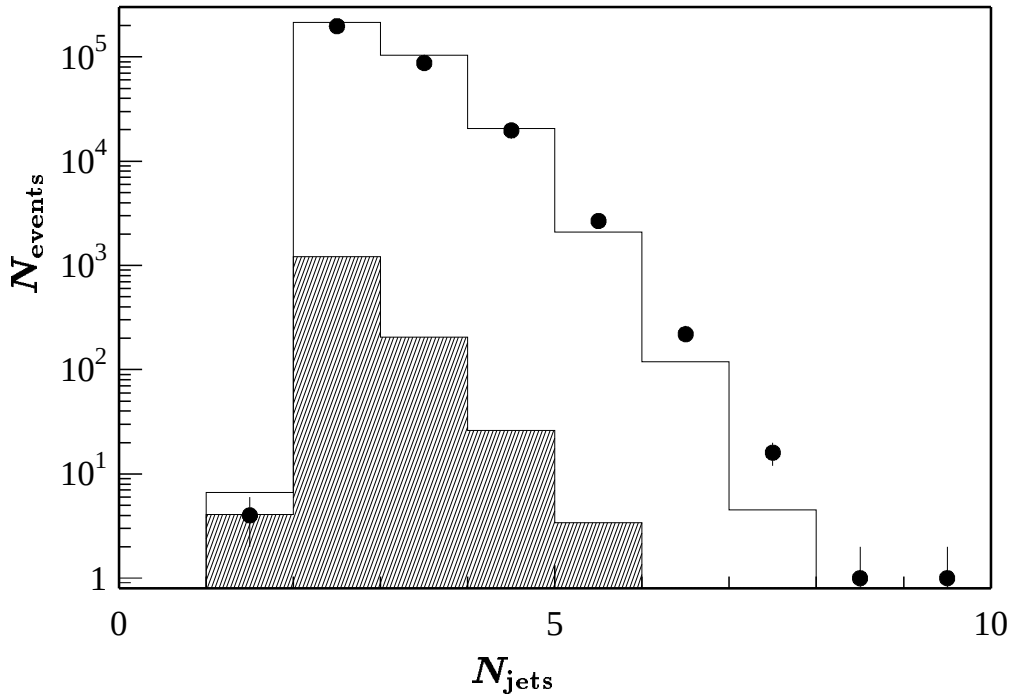


Figure 3.4: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the number of jets reconstructed per event. The hatched area denotes the background predicted by the Monte Carlo.

the above three jet types separately. Since the amount of material encountered by jets varies with the jet polar angle, this calibration is performed (and applied) in a polar angular dependent way: ten different polar angular regions are considered. For the multihadronic jet energy calibration, only events with two reconstructed jets are used. Moreover these jets are required to be balanced: their total energy should be close to $\sqrt{s}$, and their axes should be aligned, i.e. the acollinearity angle ψ defined by $\cos \psi = -\mathbf{n}_1 \cdot \mathbf{n}_2$, where $\mathbf{n}_1, \mathbf{n}_2$ are the jet axis directions, should be small. Such events are expected to be well contained in the detector; the average jet energy can therefore be constrained to the beam energy by means of the application of angle-dependent weighting factors. This procedure also improves the jet energy precision [34]. The calibration factors thus obtained are applied to all multihadron jets. Similar procedures are carried out for the electromagnetic and single hadron jet types.

For each event the so-called thrust axis is determined. The thrust axis is defined as the axis which maximises the ‘thrust’ T :

$$T = \max_{\mathbf{n}} \frac{\sum_i |\mathbf{p}_i \cdot \mathbf{n}|}{\sum_j |\mathbf{p}_j|} \quad (3.1)$$

where the summations run over all calorimetric cluster momenta p_i (considering

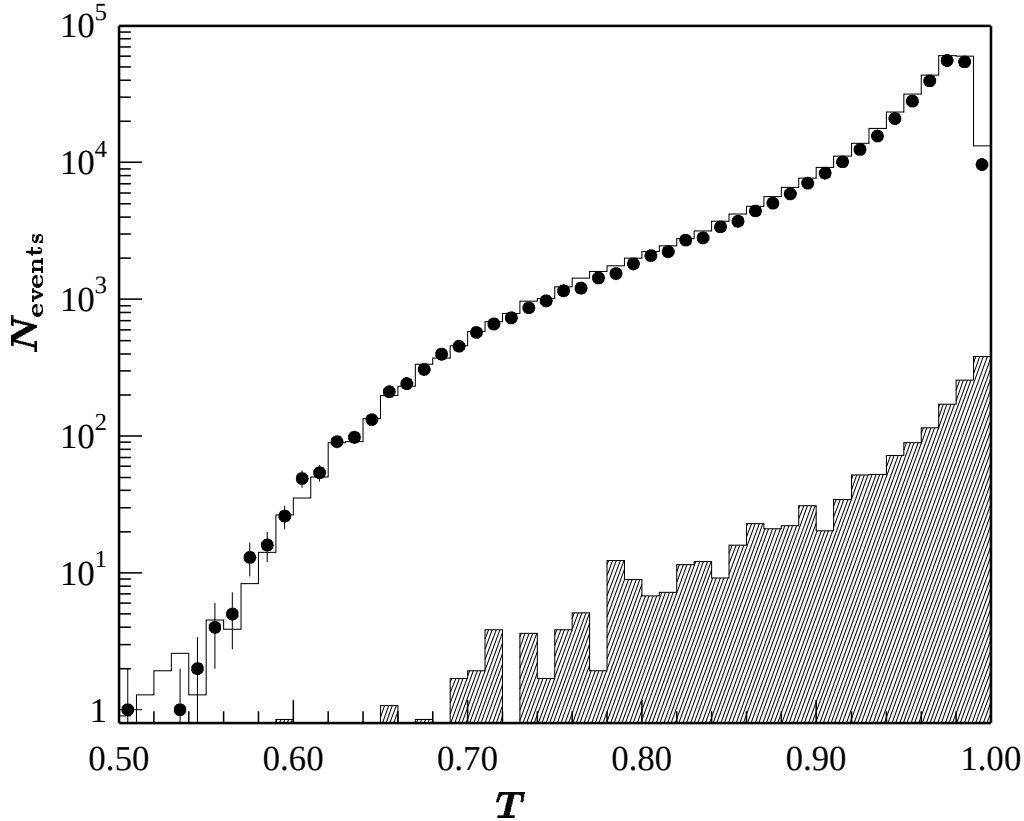


Figure 3.5: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the event thrust. The hatched area denotes the background predicted by the Monte Carlo.

the clusters to be massless particles). This procedure only yields the orientation of the thrust axis; a genuine direction is obtained by choosing the event hemisphere (with respect to the plane perpendicular to the thrust axis) containing the highest reconstructed energy. Figure 3.5 shows the thrust distribution for the 1991 data, and the corresponding Monte Carlo prediction. A very good agreement is observed. Figure 3.6 shows, for the same data and Monte Carlo samples, the polar angles of the event thrust axis. In this case too, a satisfactory agreement is observed.

The agreement between data and MC of the resulting calibrated energies is in general not perfect. For technical reasons the calibration of the energies for the data can only be performed on a subsample of the events. Inefficiencies or systematic variations in the gains of the calorimeter channels are also expected to lead to small discrepancies in the jet and total event energies. These discrepancies are corrected for in an average way: for both the data and the hadronic MC samples accepted by the selection criteria (see below) all the total event energies are scaled such that the

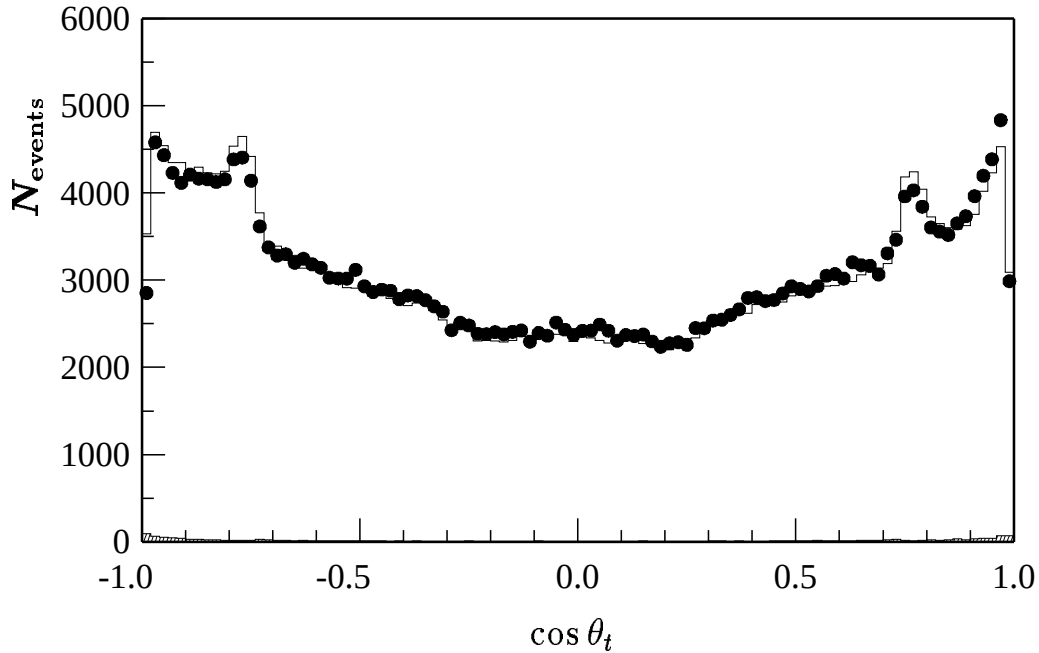


Figure 3.6: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the event thrust polar angle. The hatched area denotes the background predicted by the Monte Carlo.

observed visible distributions peak at $E_{\text{vis}}/\sqrt{s} = 1$ (see section 3.2.4). In the MC case, these factors are also applied to the background MC samples. In all cases, the resulting scale factors differ from one by less than 2%.

3.2 Selection of Hadronic Events

3.2.1 Event signatures

The signatures of leptonic events are well known. Electrons do not decay, and muons hardly ever do inside the detector. The τ is a short-lived particle, but its decay modes have been analysed thoroughly.

The situation is less clear in the case of hadronic events. It is rather well established that the particles emerging from a primary annihilation are the quarks, anti-quarks and gluons (commonly denoted as partons). However, a distinctive experimental feature associated with QCD is the fact that only colour singlets can be observed freely in nature. This means that the (coloured) quarks and gluons have to transform into the observable colourless hadrons. This process – called fragmentation – takes place at relatively large ($\gtrsim 1$ fm) interaction distances, corresponding to low momentum transfers. In this regime the QCD running coupling constant $\alpha_s(Q^2)$ becomes large. This implies that fragmentation is essentially a non-perturbative process; as a result, no description of fragmentation starting from

first principles (i.e. a QCD Lagrangian) exists at this moment. Instead, several phenomenological fragmentation models have been developed to translate as much as possible of the experimental knowledge of the physics processes into convenient algorithms.

The origin of these models is the assumption that for large distances, the interquark potential rises linearly with distance. As the primary quarks and gluons separate, the potential energy needed for the parton binding at some point becomes large enough for it to be energetically favourable to create another $q\bar{q}$ pair. Mesons are then formed by binding the primary quark to an antiquark resulting from this break-up. This process continues until some cutoff energy is reached.

By itself, this picture does not provide a mechanism to account for baryon production. This is included in an *ad hoc* way, through the assumption that a certain fraction of the time, the breakup of the quark bond proceeds through the creation of diquark – anti-diquark pairs; the binding of a primary quark to a diquark then produces a baryon.

The fragmentation functions specifying the momentum distribution of the hadrons arising from a fragmentation process are tuned to existing data and depend both on the quark flavour and on the hadrons involved. Also the description of the relative hadron production rates and the decay processes depends largely on phenomenological input.

Nevertheless, several Monte Carlo programs based on models of this kind have been remarkably successful in describing the experimental data [43]. The performance of these programs plays an essential role in attaining the precision aimed for in the high-statistics hadronic channels of the LEP experiments.

The qualitative features of hadronic events at LEP energies can be deduced from the above picture. First, only hadrons are produced in the fragmentation process; leptons and neutrino's only follow from the (semi-)leptonic decays of hadrons. Since (semi-)leptonic branching fractions are generally small this implies that on average, little energy will escape detection in the form of neutrino's. This in turn means that, provided the detector offers a good hermeticity (99% of 4π in the case of the L_3 detector), the total energy measured by the detector (the 'visible energy') will be close to $\sqrt{s}$ and the visible momentum will be well balanced. Second, the fragmentation process will continue until the energy has degraded far enough for it to become energetically impossible to create more particles. This means that the higher $\sqrt{s}$, the higher the average particle multiplicity will be. Given the fact that most particles produced in the fragmentation are pions, isospin symmetry arguments will then also require an increase of the charged particle multiplicity with increasing $\sqrt{s}$. At LEP energies, the average charged particle multiplicity is found to be approximately 21 [44]. Figure 3.7 shows a typical hadronic event as seen in the L_3 detector.

In practice, the generation of events by hadron Monte Carlo programs proceeds in four steps:

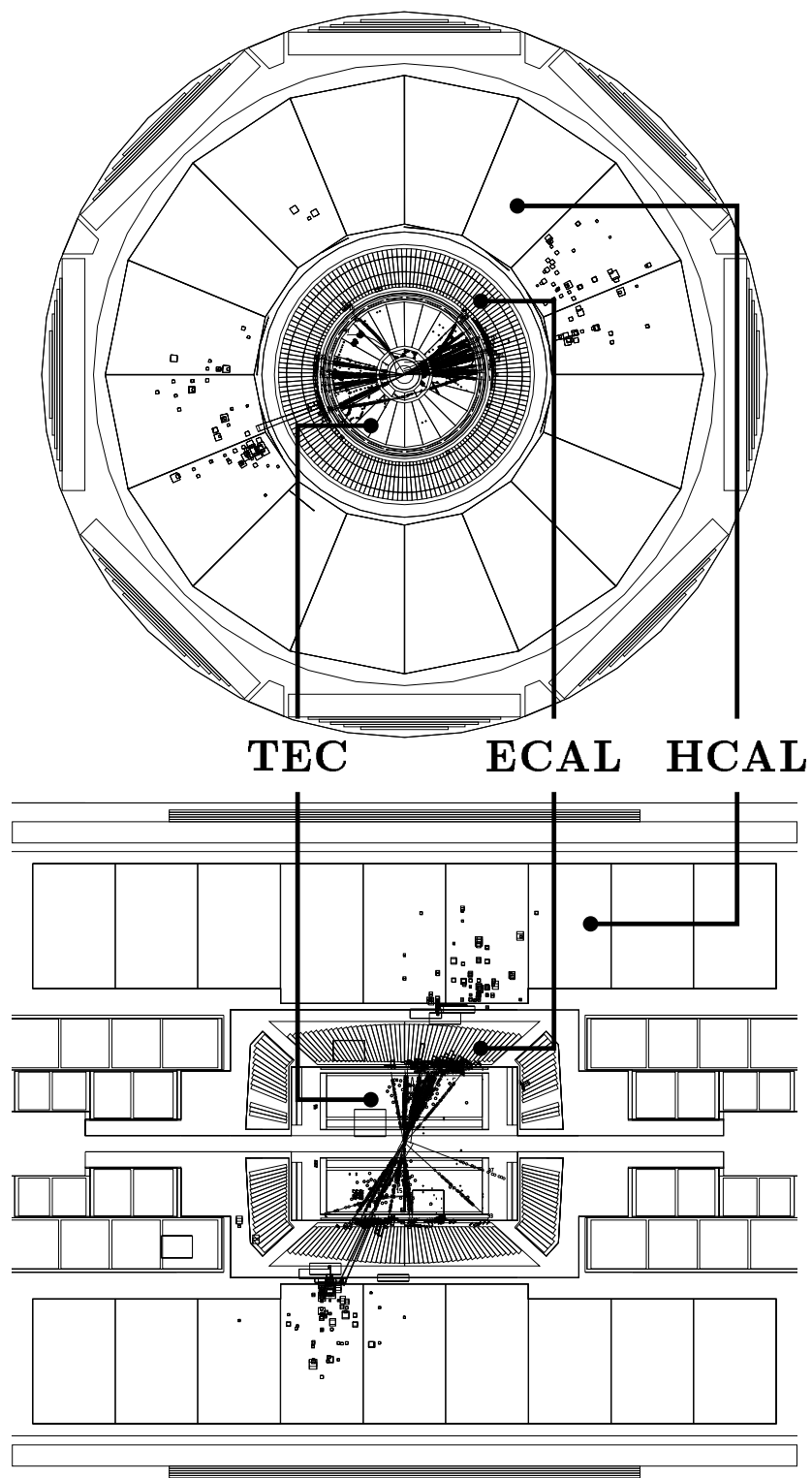


Figure 3.7: *A typical hadronic 2-jet event as seen in the L_3 detector.*

1. *The primary annihilation.* It is this phase which is sensitive to electroweak effects. To the lowest order process one may or may not apply electroweak and photonic corrections.
2. *QCD corrections to the primary annihilation.* Due to the fact that α_s is relatively large, corrections to the lowest order $e^+e^- \rightarrow q\bar{q}$ process have to be taken into account. The so-called Matrix Element (ME) approach consists of an order-by-order Feynman diagram calculation which should in principle give the correct results. This calculation has however only been performed up to $\mathcal{O}(\alpha_s^2)$. Thus in the ME approach, only final states with up to four partons can be taken into account. Given the fact that the corrections to be applied are sizable, this approach cannot be expected to describe the full perturbative branching process.

For this reason, one often reverts to the so-called Parton Shower (PS) algorithms. These algorithms describe the exact matrix elements in so-called Leading Logarithmic Approximation (LLA)[45]. They are well-suited to a probabilistic treatment, and hence to applications in MC programs. Additional (nonleading) interference effects can either be accommodated in the branching algorithm or be put in by hand. The perturbative branching process stops when the momentum transfer scale Q^2 becomes so small and the corresponding coupling constant too large for a perturbative treatment to remain valid; a cutoff value of several GeV is typical.

3. *Fragmentation.* Two conceptually different models are nowadays used to describe the fragmentation process:
 - String fragmentation. In this approach the interquark potential is viewed as a ‘colour flux tube’ or *string*. The radiation of gluons introduces ‘kinks’ in these strings. Strings break apart (usually through the formation of a $q\bar{q}$ pair) if the ‘string tension’ (the energy density) becomes too high. This algorithm can be applied to either the ME or the PS approach.
 - Cluster fragmentation. This algorithm starts by splitting all gluons generated by the perturbative branching into $q\bar{q}$ pairs. Quarks and antiquarks which are adjacent in momentum space are combined into (colourless) hadron clusters; high mass clusters are split into smaller clusters. These clusters subsequently decay into the final state hadrons. Since (apart from gluon splitting) no branching takes place, this algorithm can only be applied in conjunction with the PS approach.
4. *Decay processes.* At this stage, there is no more use for specific models. Instead, one relies on experimental information on particle decays. For the data analysed in this thesis, there is no strong dependence on the detailed description of these processes.

The calorimetric coverage of the L_3 detector extends to polar angles as low as 5° , whereas full sampling of charged tracks stops at 45° . Given the resulting higher efficiency for calorimetric measurements, a calorimetric selection clearly offers the best possibilities for a precision measurement. The data in this thesis are therefore only based on calorimetric selection.

3.2.2 Monte Carlo predictions

All ‘physics’ processes (i.e. processes caused by e^+e^- interactions) expected to contribute to the observed signal were simulated. One also expects beam-gas events to play a role, but unfortunately no Monte Carlo generator program exists capable of simulating such processes. This problem will be addressed in chapter 4.

The most important source of Monte Carlo information are the simulated hadronic events themselves. For this purpose the JETSET [46] generator program was used. This program features a full Standard Model description of the lowest order scattering process and first order (unexponentiated) bremsstrahlung corrections. The program has both a ME and a PS option; the Monte Carlo sample used in the present analysis was created using the PS option. Gluon interference effects causing corrections to the (LLA) branchings are introduced as additional constraints. The fragmentation process is described using the string fragmentation approach. This program is only used to calculate the detector acceptance for hadronic final states. Since bremsstrahlung photons usually have a small energy, the absence of a more complete treatment of bremsstrahlung corrections (cf. multiple hard photon emission) is not expected to cause large inaccuracies. The events were generated at a center-of-mass energy of 91.25 GeV, corresponding to the Z resonance peak where most of the data were taken.

The KORALZ [47] program was used for the estimation of the background from the $e^+e^- \rightarrow \tau\tau(\gamma)$ process. This program takes into account higher order initial state and first order final state bremsstrahlung corrections, $\mathcal{O}(\alpha)$ electroweak corrections and polarisation effects. Experimental data are used to quantify the relative τ decay rates.

In principle one also has to take into account the background from the process $e^+e^- \rightarrow \mu^+\mu^-(\gamma)$. However, from a Monte Carlo study (using the same KORALZ program) it was concluded that any background from this final state can be neglected altogether.

In order to estimate the $e^+e^- \rightarrow e^+e^-(\gamma)$ background contamination, the BABAMC program was used. This program is described in somewhat more detail in chapter 4. Here this program is used with a polar angle cutoff in order to avoid the forward elastic peak (see chapter 1). Events were generated in the region $6^\circ < \theta < 174^\circ$, corresponding to a peak cross section of 13.42 nb.

Lastly, the DIAG36 [48] program was used to study the background contamination from the two-photon process $e^+e^- \rightarrow e^+e^-q\bar{q}$. This program evaluates the low-

est order QED contributions to four-fermion final states, of which only the hadronic final states were used. Fragmentation and decay processes were simulated using the JETSET program. The cross section resulting from the program is modified by hand [49], the reason being that the program only gives the contributions from ‘pointlike’ processes. Such a treatment is insufficient for the present analysis: for low momentum transfers Q_1^2, Q_2^2 in the two-photon channel proper (see figure 3.8) nonpointlike effects become significant. The approach taken is to extrapolate lower energy measurements [50] to LEP energies using the so-called Generalised Vector meson Dominance Model (GVDM) [51]. This model assumes that at low Q_1^2 and/or Q_2^2 , the photons effectively behave as vector mesons which subsequently collide. To estimate the acceptance for this process, a cut was imposed on the generator level (i.e. before detector simulation): only final state particles emerging from the interaction point under angles of at least 4° with respect to the beam axis and having a summed energy of at least 10 GeV were considered. With this cut, the generated cross section at $\sqrt{s} = 91.25$ GeV is approximately 5.5 nb. However, due to the above mentioned extrapolation, this cross section is presumably only accurate to within a factor of two [49].

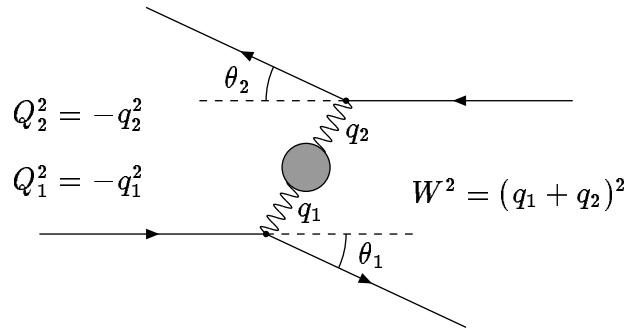


Figure 3.8: *Kinematic variables in two-photon scattering.*

3.2.3 Preselection

Before describing the final selection criteria, it is important to consider once more the primary data samples used. The total amount of data accepted by the L₃ triggers is so large as to render the analysis of the complete data sample prohibitive. For this reason a preselection of data is required. Various physics event selection algorithms are applied to the incoming data. Most of the selection algorithms (notably the $q\bar{q}$ selection) create a separate ‘data stream’; such a stream then contains all events accepted by the preselection in question and serves as the reference sample for the

corresponding analysis. A master reference sample contains all events accepted by any of the physics preselections.

In the case of the $q\bar{q}$ preselection, the algorithm consists essentially of a loose form of the final selection criteria, backed up however by two additional requirements. The first one is a rather loose cut on the number of scintillators fired within 15 ns (corrected for time-of-flight) of the bunch crossing time: for the 1990 (1991) data taking period at least 6 (7) hits were required. The second backup requirement asks for at least four charged tracks with $p_T > 100$ MeV in the CTD. As verified with the help of Monte Carlo event samples and small data samples which were visually inspected and identified, the logical OR of these selections is almost fully efficient; the purity of the resulting data stream is approximately 75%.

In the remainder of this chapter, it will always be assumed (unless stated otherwise) that all data have been subjected to the preselection filter.

3.2.4 Event selection criteria

The event reconstruction quantities which are used in the present analysis are:

- The visible energy E_{vis} . This quantity is obtained by adding the (calibrated) energies of all the jets and all the remaining ('unused') clusters. In the analysis this sum is always normalised to $\sqrt{s}$ (on a fill-by-fill basis) in order to avoid introducing an implicit energy dependence.
- The longitudinal and transverse energy imbalances $E_{\parallel}$ and E_T . These quantities are obtained by constructing the vectorial sum of all cluster momenta $E_i \hat{\mathbf{p}}_i$ (neglecting all particle masses):

$$\mathbf{E} = \sum_i E_i \hat{\mathbf{p}}_i. \quad (3.2)$$

The imbalances are then defined as

$$E_{\parallel} = |E_z|/E_{\text{SRC}}, \quad E_T = \sqrt{E_x^2 + E_y^2}/E_{\text{SRC}} \quad (3.3)$$

where E_{SRC} is defined as the simple energy sum of all the calorimeter clusters. The reason for using the quantity E_{SRC} is that it allows a consistent treatment of the cluster momenta. The previously defined E_{vis} gives the best energy precision.

- The cluster multiplicity. This is simply the number of reconstructed SRCs, with the additional requirement that the SRCs contain at least one of the following objects:
 - a HCAL cluster;
 - an ECAL bump consisting of at least two crystals;

- a charged track measured in the TEC.

This requirement suppresses random noise clusters in the ECAL.

The granularity of the ECAL is much higher than that of the HCAL. Thus, for the same number of incident particles the reconstructed cluster multiplicity depends on the presence of the full ECAL. In the 1990 data taking period, the ECAL endcaps had not yet been installed. For this period, the number of reconstructed clusters in the small angle region therefore differs significantly from the corresponding number in the large angle region. This is most relevant for high multiplicity events, where (due to the better granularity of the ECAL) nearby energy depositions are more easily resolved than in the HCAL. In order to take this difference into account, the requirement on the cluster multiplicity is made to depend on whether the thrust axis lies inside or outside the barrel region. For this purpose, an event for which the polar angle of the thrust axis satisfies $|\cos \theta_t| < 0.743$ is defined as a ‘barrel’ event; otherwise it is said to be an ‘endcap’ event.

Figure 3.9 shows the Monte Carlo expectations for the two most important quantities used in the event selection, the visible energy (scaled by $\sqrt{s}$) and the cluster multiplicity. This figure clearly confirms the picture sketched above; it also indicates approximately where the cuts should be placed.

The event selection criteria, based on the observations made above, are then as follows:

1. For barrel events the number of clusters has to be at least either 11 (for the 1990 period) or 13 (for the 1991 period). Figure 3.10 shows the corresponding distributions². As can be seen from the figure, this cut distinguishes rather sharply between low multiplicity e^+e^- and $\tau\tau$ final states on one hand and high multiplicity hadronic final states on the other hand. The discrepancy seen at the high end of the distribution for the 1991 data is partially understood to be due to an incorrect hadronic cross section used in the simulation of the interactions of charged pions with the detector material. Any consequence of this error on the low end of the distribution is expected to be covered by the approach discussed in section 4.2.6.
2. For endcap events the number of clusters has to be at least either 8 (for the 1990 period) or 17 (for the 1991 period). The distributions are again shown in figure 3.10. This requirement is completely equivalent to the one in the barrel region. However, in this case a much more significant Bhabha background is accepted, for two reasons:
 - The Bhabha cross section in the endcap region is dominated by t -channel photon exchange singularity and much larger than in the barrel region;

²The normalisation of the Monte Carlo histograms in this section is an absolute one, based on the cross sections for the various processes.

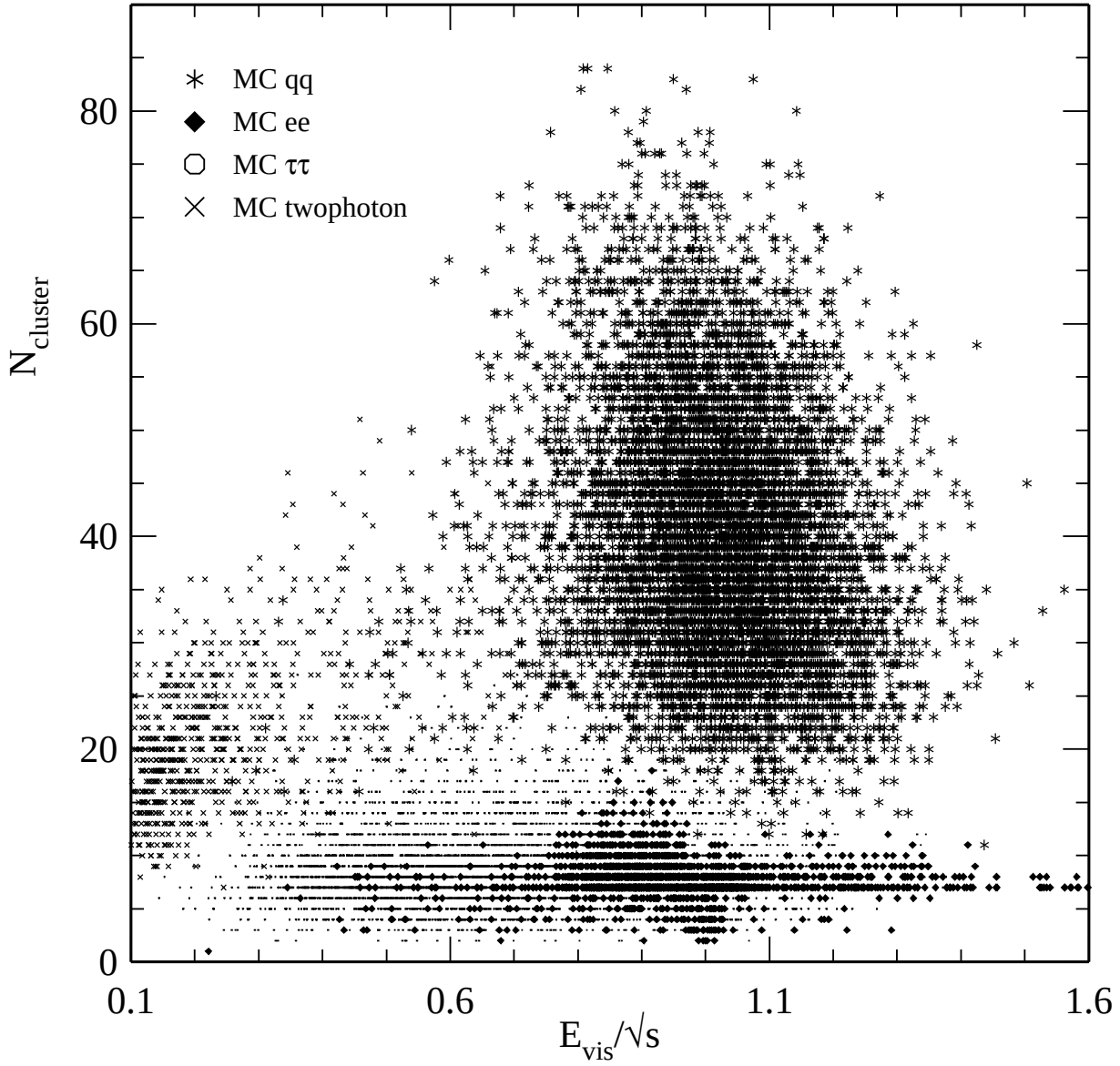


Figure 3.9: Monte Carlo signal and background predictions for the event multiplicity and for the visible energy (1991 detector setup).

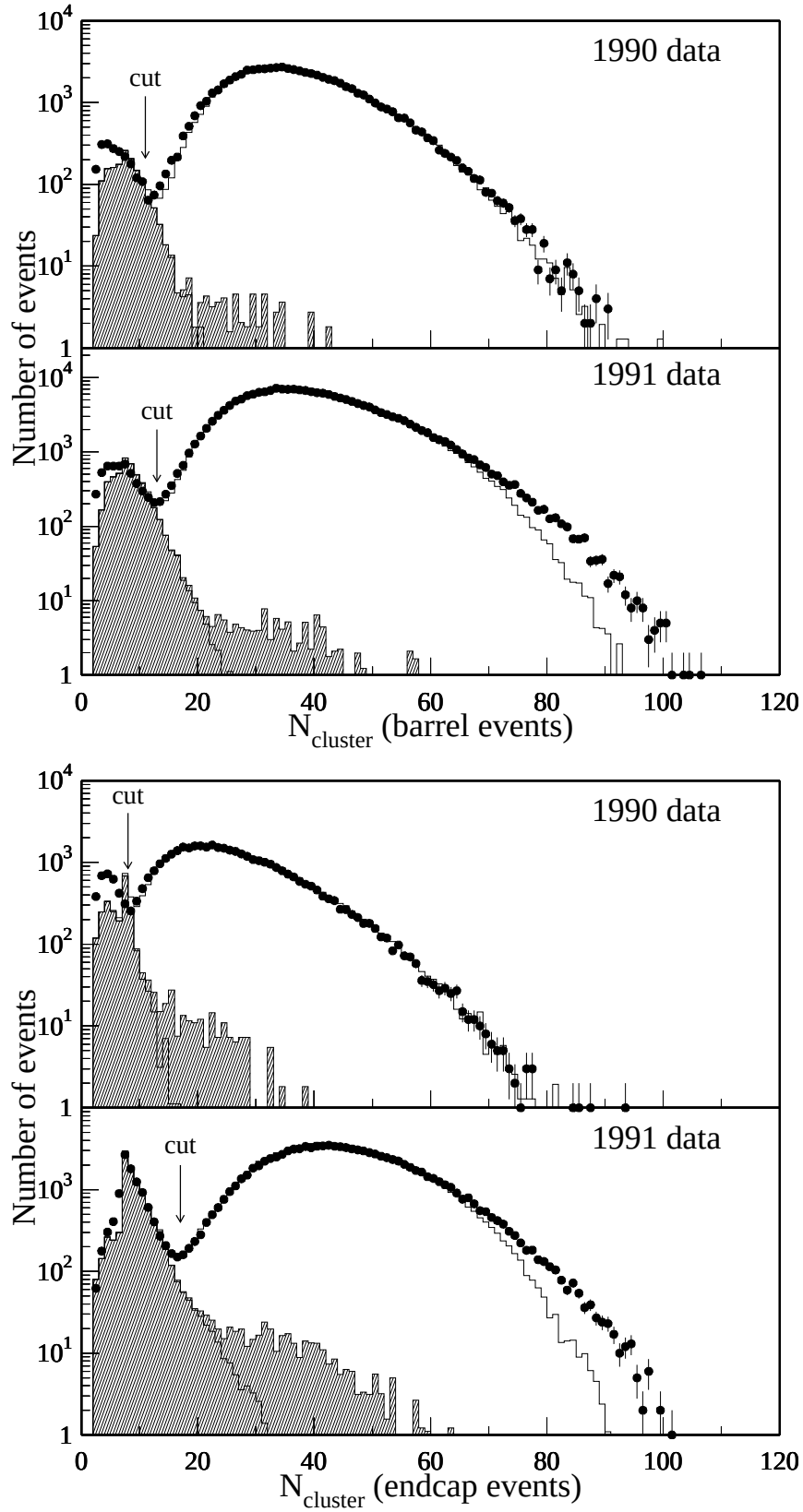


Figure 3.10: Cluster multiplicity distributions for the 1990 and 1991 data, compared with the corresponding Monte Carlo predictions. Top: barrel events. Bottom: endcap events. The hatched area represents the contribution expected from background processes; the second line inside this area denotes the same contribution, but with the two-photon contribution omitted.

- Due to the absence of the ECAL endcaps in the 1990 data taking period, calorimetry was performed with the HCAL only. This leads to much larger fluctuations, in particular in the number of reconstructed SRCs.
3. The visible energy has to lie between $0.5\sqrt{s}$ and $1.5\sqrt{s}$. The distributions are shown in figure 3.11. As is clear from figure 3.9 already, the minimum energy requirement rejects most of the two-photon background and a fair fraction of the tau background. For the 1990 data, the cut on the high end of the

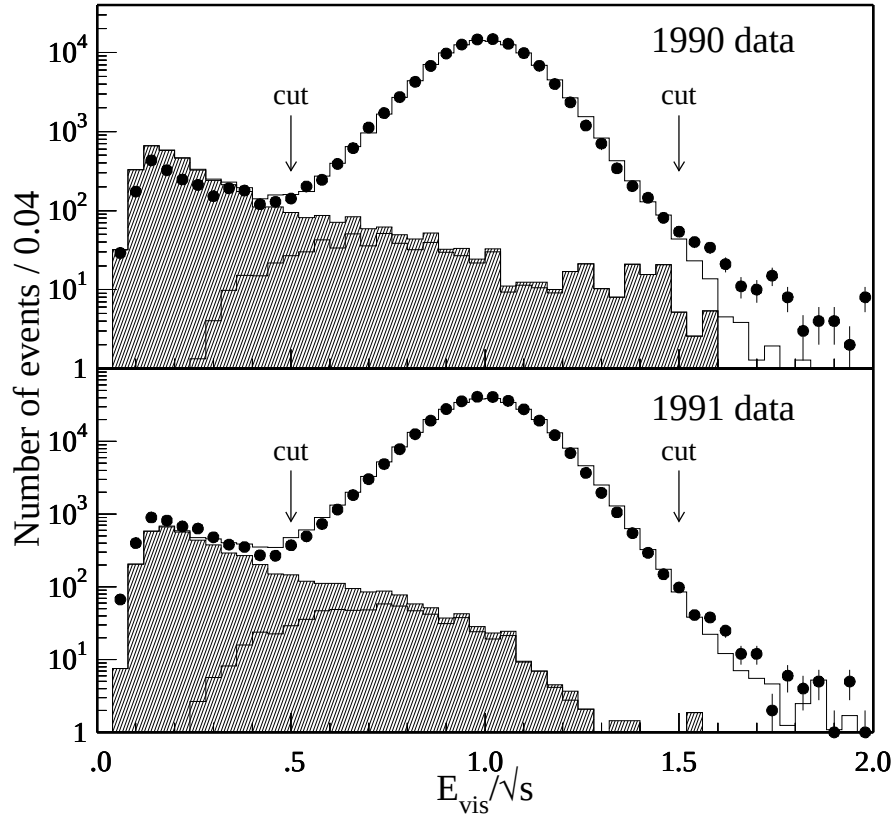


Figure 3.11: Scaled visible energy distributions for the 1990 and 1991 data, compared with the corresponding Monte Carlo predictions. The hatched area represents the contribution expected from background processes; the second line inside this area denotes the same contribution, but with the two-photon contribution omitted.

distribution mainly removes Bhabha events. The presence of this Bhabha background in the 1990 data is due to the absence of the ECAL endcaps. Without these endcaps the separation in multiplicity between hadronic and Bhabha events is less pronounced and the corresponding multiplicity cut less

selective. For reasons of consistency the same cut is applied for the 1991 data taking period.

4. The transverse energy imbalance has to be smaller than 0.5. The distributions are shown in figure 3.12. The main purpose of this cut is to remove noise background. A visual inspection of the events in the high end tail for the 1991 data taking period leads to the conclusion that the observed excess is due to cosmic muon events, interactions of the $e^\pm$ beams with the wall of the vacuum pipe and correlated electronic noise caused by support tube movements. Similarly, the excess observed for the 1990 data is expected to be due to beam-wall and cosmic muon background.
5. For the 1990 period, the longitudinal energy imbalance has to be smaller than 0.5. For the 1991 period this cut is somewhat relaxed to 0.6. The distributions are again shown in figure 3.12. The observed agreement between data and MC is good. However, relaxing this criterion would bring the cut value quite close to the corresponding preselection cut (0.55 and 0.75 for the 1990 and 1991 data respectively). For this reason it is necessary to keep the above values (see also the discussion in section 4.2.4).

Each of the distributions shown in figures 3.10-3.12 is made after application of all other cuts. Thus, all plots show the events rejected by the selection criterion in question.

One additional cut is applied to remove events caused by correlated noise in the HCAL:

6. The hit multiplicity observed in the HCAL (more precisely, the number of towers which are either associated with HCAL clusters or have an energy of at least 1 GeV) is required to be at most 800.

As shown in figure 3.13 the effect of this cut is small. The discrepancies remaining inside the acceptance region are presumably due to inaccuracies in the HCAL calibration. However, the loose cut applied ensures that this has no effect on the total cross sections. The plot shown is for the 1991 data; the same conclusion holds for the 1990 data sample.

3.2.5 Corrections to Monte Carlo events

The simulation of Monte Carlo events generally provides a fairly good description of all physics process taking place. However, some imperfections relevant for the present analysis remain. For this reason, *a posteriori* corrections are applied to the Monte Carlo events:

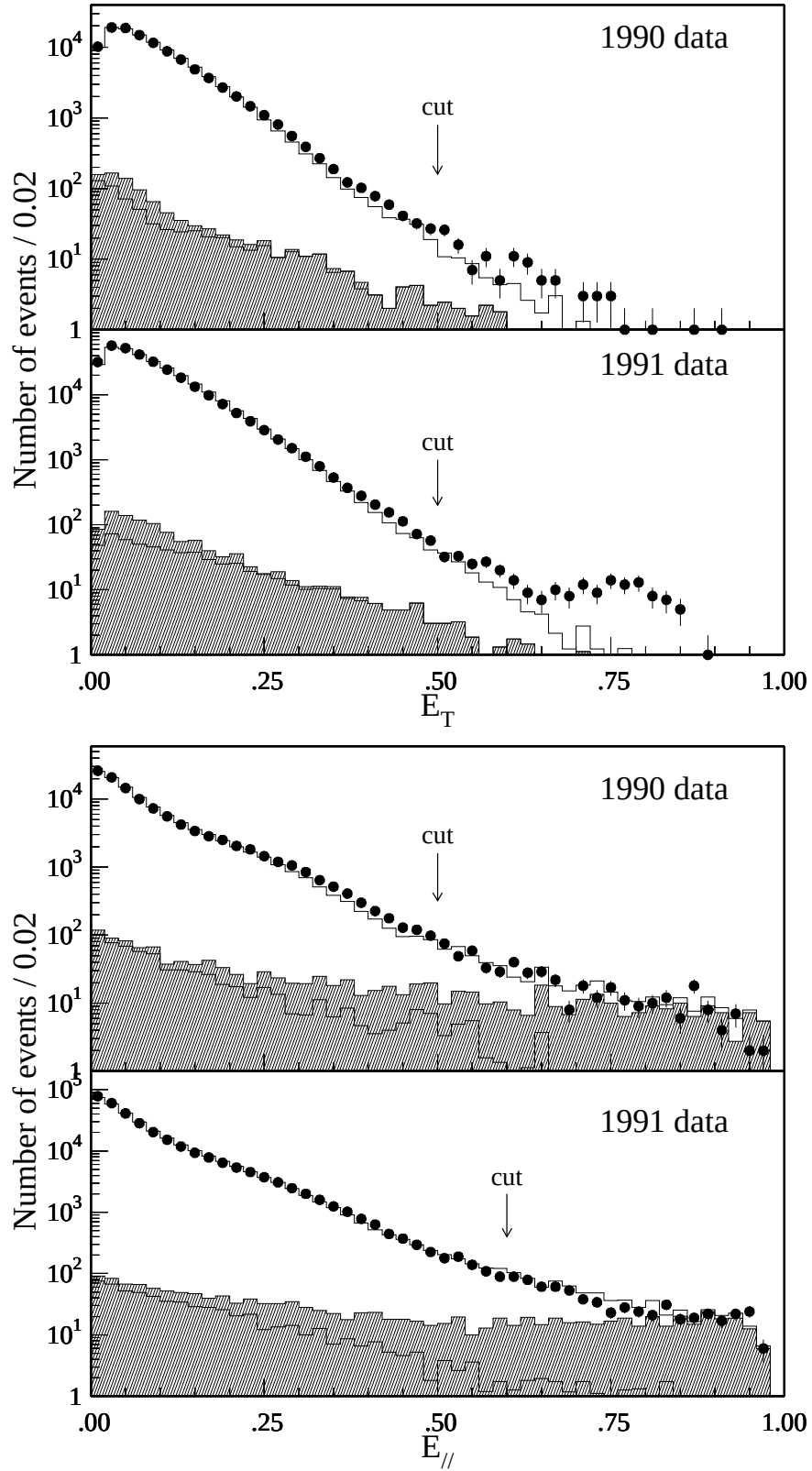


Figure 3.12: Transverse (top figure) and longitudinal (bottom figure) energy imbalance distributions for the 1990 and 1991 data, compared with the corresponding Monte Carlo predictions. The hatched area represents the contribution expected from background processes; the second line inside this area denotes the same contribution, but with the two-photon contribution omitted.

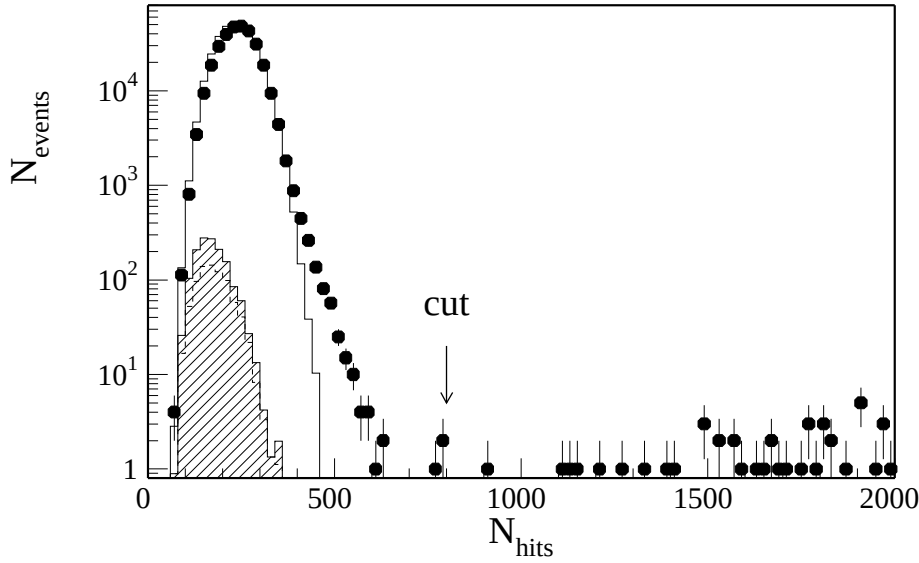


Figure 3.13: Comparison of observed hit multiplicity in the HCAL with Monte Carlo predictions (1991 data).

- *corrections due to uranium noise:* the radioactivity of the absorber material of the HCAL not only offers a calibration method, it also represents a continuous source of background. As was mentioned in subsection 3.1.3, most of the background from the uranium signal is suppressed by a simple energy requirement on the individual hits; however, when uranium noise coincides with the signal from a physics interaction the two are added. This fact has to be taken into account at the stage of the event simulation; a random uranium noise spectrum is added to the MC events on the hit level.
- *corrections due to random background:* other sources of random signals, such as electronic noise or overlapping events ('pile-up') are also present. Assuming an event rate of 1 Hz and a bunch crossing rate of 45 kHz one expects pile-up at a rate of $\sim 10^{-5}$. The effect on the measured energy is expected to be negligible and is disregarded in this analysis; however, a slight increase of the cluster multiplicity is expected which is particularly significant for low multiplicity events. This effect can be studied by considering the so-called 'beam-gate' data stream. Beam-gate events are events taken at beam crossing time without any particular trigger requirement. Thus they represent an unbiased source of background. For the present analysis, a distribution of reconstructed numbers of beam-gate clusters (as found in this 'beam-gate' stream) is added to the Monte Carlo distributions. For the 1990 and 1991 data taking periods, the mean numbers of beam-gate clusters are found to be 0.55 and 0.77 respectively.

It should be noted that this treatment is not completely bias free; this problem will be addressed in section 4.2.6.

- *corrections due to the simulation energy cutoff*: the GEANT package applies an energy cutoff on neutral hadrons of 10 keV. Hadrons with kinetic energies below this value are not tracked. As a consequence, slow neutrons produced in the interactions between the particles produced in the primary interaction and the uranium of the HCAL are not treated properly. Since these neutrons may interact in the scintillator counters within the time gate of 15 ns, a discrepancy (between data and MC) is observed in the distribution of the number of fired counters. Although this quantity is not used in the final event selection, it is relevant in the preselection as it affects the estimated amount of background. This is corrected for on an *ad hoc* basis: for the $\tau\tau(\gamma)$ and $ee(\gamma)$ final states a preselection exists just as for the $q\bar{q}$ final states. First, the same selection criteria are applied to the MC data as those needed to obtain a relatively pure subsample in the corresponding real data. Subsequently the MC are ‘smeared’ to resemble the data by adding a contribution of randomly fired scintillator counters. This contribution is chosen such that a reasonable agreement is obtained between data and MC in the mean number of fired scintillators as a function of the event polar angle and the cluster multiplicity. An example of the changes in the distributions for the 1991 Bhabha event sample thus obtained is given in figure 3.14. One observes a good agreement between data and Monte Carlo distributions. The 1991 data passed through the selection criteria discussed in section 3.2.4 are compared to the corresponding Monte Carlo prediction in figure 3.15. Here one observes that some discrepancies between these two still remain. The Monte Carlo sample is however dominated by the $q\bar{q}$ contribution, for which the preselection based on the scintillator multiplicity is of minor importance only; hence the $q\bar{q}$ selection efficiency is not affected by these discrepancies. It should be stressed that this smearing is only relevant for endcap events; it mainly affects the Monte Carlo predictions at the low end of the distribution of the number of clusters (see figure 3.10).

In order to treat real and simulated data consistently, the above corrections are applied to the Monte Carlo events before they are subjected to the preselection and final selection criteria.

3.2.6 Selection results

The above-mentioned selection criteria are applied both to the complete 1990 and 1991 data and to the Monte Carlo samples. However, in order to obtain reliable cross section estimates one also has to take into account the fact that at times, the performance of the detector becomes too degraded for the data taken during these

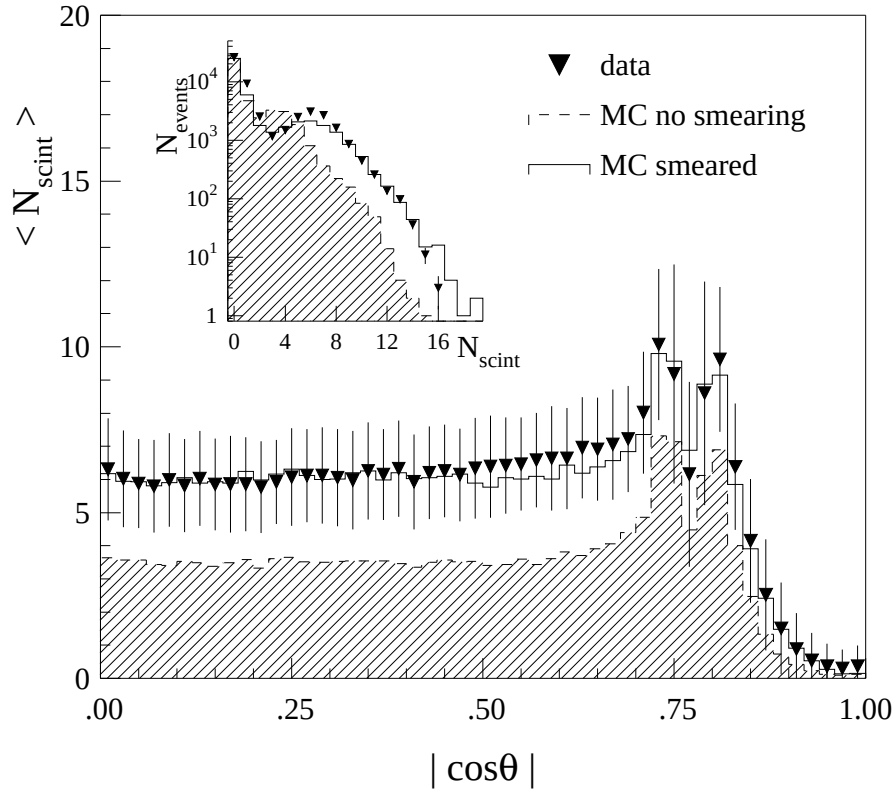


Figure 3.14: Comparison between 1991 Bhabha events and corresponding Monte Carlo predictions (before and after corrections) of the average number of fired scintillator counters as function of the event polar angle, and of the overall distribution of the number of fired counters. The error bars in the main plot indicate the spread around the mean.

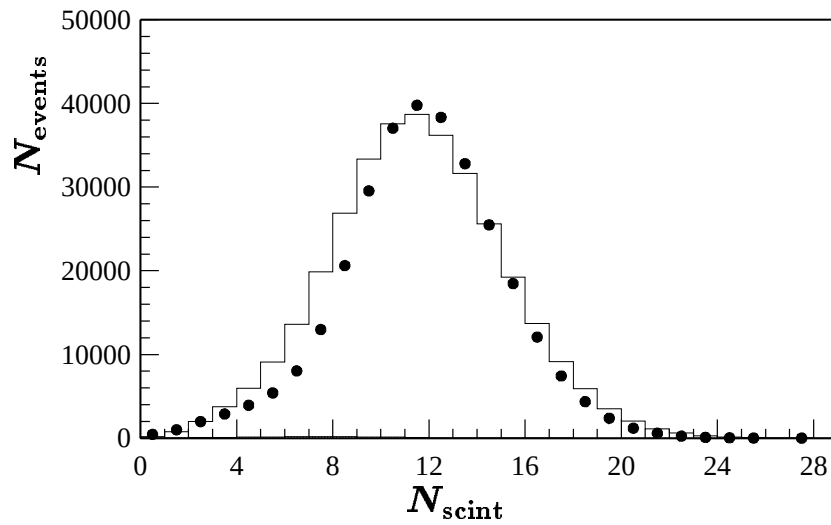


Figure 3.15: Comparison between 1991 hadronic data and corresponding Monte Carlo predictions, of the number of fired scintillators.

periods to be useful for analysis. For the purpose of the present analysis, runs were marked as ‘bad’ if one or more of the following conditions were fulfilled:

- One or more BGO ‘rings’ dropped from the readout during the run; this is flagged by the absence of a signal from the unparsed crystal in the ring. It is technically impossible to simulate this source of inefficiency in the Monte Carlo events;
- Either six or more BGO rings, or one or more HCAL modules, were inactive in a run;
- A run was ‘noisy’. This noise can be caused either by a specific component of the readout system itself or by external factors. The first cause is registered in the subdetector databases; the second one must be identified during the data analysis;
- The energy trigger (including the luminosity trigger) experienced problems, usually because of excessive background rates.

The sizes of the resulting ‘good run’ samples used for this analysis are summarised in table 3.1. Since part of these deficiencies (missing BGO rings) are introduced in the Monte Carlo events as well, the runs classified as ‘bad’ are excluded both for the real and simulated data.

	1990		1991	
	N_{events}	ϵ (%)	N_{events}	ϵ (%)
MC $q\bar{q}$	169232	98.64 ± 0.03	531265	99.23 ± 0.01
MC e^+e^-	26130	0.34 ± 0.04	48117	0.01 ± 0.01
MC $\tau^+\tau^-$	23743	7.01 ± 0.17	89729	4.00 ± 0.07
MC two-photon	30495	1.03 ± 0.05	59888	1.28 ± 0.05
Data	187124		418934	

Table 3.1: *Samples used for the analysis of the 1990 and 1991 data. The numbers for the Monte Carlo samples are before application of the preselection.*

Table 3.1 also gives the computed selection efficiencies for each Monte Carlo sample. The errors quoted on these efficiencies merely reflect the statistical uncertainty following from the finite size of these MC samples. They are calculated using the binomial formula

$$\Delta\epsilon = \sqrt{\epsilon(1 - \epsilon)/N_{\text{events}}} \quad (3.4)$$

where N_{events} is the number of generated Monte Carlo events. The quoted efficiencies do not take into account the time dependent detector-related effects which are applied in the analysis. When computing cross sections in the next chapter, these effects are properly taken into account.

Chapter 4

Hadronic Cross Sections

The hadronic lineshape measurement involves the determination, as a function of the center-of-mass energy $\sqrt{s}$, of the number of events in which a $q\bar{q}$ pair couples to an intermediate photon or Z boson. Generally, the conversion between numbers of events and cross sections proceeds through a normalisation to the number of colliding particles. The quantity of interest is the so-called (time-) integrated luminosity:

$$\sigma = N_{\text{events}} / \int \mathcal{L} dt \quad (4.1)$$

In principle, for head-on collisions the instantaneous luminosity $\mathcal{L}$ can be expressed as

$$\mathcal{L} = \frac{N_{e^+} N_{e^-} n_b f_{rev}}{4\pi \sigma_x^* \sigma_y^*} \quad (4.2)$$

where it is assumed that the numbers of electrons and positrons $N_{e^\pm}$ are evenly divided among the n_b bunches per beam [15]; σ_x^* (σ_y^*) is the rms spread of the beam in the x (y) direction; their product represents the geometrical cross section of the bunches at the interaction point. Evidently a luminosity determination based on the measurement of the beam parameters occurring in formula 4.2 is difficult. One therefore normally relies on the fact that the resulting normalisation is the same for all physical processes. One thus measures the number of events for a process with a known cross section; eq. 4.1 then yields the integrated luminosity. The cross section for this process should be large in order to reduce statistical fluctuations. The small angle Bhabha scattering is such a process. The first part of this chapter is devoted to the measurement of the L_3 integrated luminosity. It should be stressed that the instantaneous luminosity is of no relevance for the derivation of the results in this thesis; henceforth the integrated luminosity will be denoted by $\mathcal{L}$.

The selection of hadronic events is described in chapter 3. Combining the results of this event selection with those of the luminosity measurement the hadronic cross sections obtained for the 1990 and 1991 data taking periods are evaluated. The simple formula 4.1 is only valid for the case where the selection of events is fully

efficient and free from background; to arrive at a final value cuts and corrections are necessary. Finally various consistency checks can be applied to the obtained results; the outcome of these checks indicates the importance of systematic uncertainties.

4.1 Luminosity Measurement

The physics process that is invariably chosen for the luminosity measurement in e^+e^- experiments is Bhabha scattering $e^+e^- \rightarrow e^+e^-(\gamma)$. As stated in chapter 1, at small angles this process is dominated by t -channel photon exchange; this is a QED process and well understood. Thus, to a good approximation the luminosity measurement can be regarded as being independent of the Standard Model parameters, the determination of which is the subject of this thesis.

The luminosity analysis in the L_3 collaboration is performed using the information from the small-angle BGO calorimeters (see chapter 2); no attempt is made to distinguish electrons and photons.

4.1.1 Luminosity trigger

As in the case of the trigger signals derived from the barrel and endcap ECAL, the luminosity trigger uses the (digitised) analog sums of the 19 crystal signals belonging to the same azimuthal sector. The resulting 2×16 signals are subject to the following trigger conditions [37]:

- *back-to-back trigger*: at least 15 GeV is measured in both calorimeters, with an azimuthal difference of at most two sectors;
- *asymmetric double-tag trigger*: at least 25 GeV is measured in one calorimeter, and at least 5 GeV in the other;
- *single-tag trigger*: at least 30 GeV is measured in one of the calorimeters.

The single-tag trigger is merely used to estimate the efficiency of the other two triggers and is ‘prescaled’ by a factor of 40 (i.e., only every 40-th trigger is an event actually written onto tape). From the accepted luminosity events one can derive an overall trigger efficiency; from this it is concluded that the trigger can be regarded as fully efficient. In most of the 1991 data taking period, the luminosity trigger was prescaled by a factor of two in order to reduce the event rate (and thereby the dead time). This does not significantly degrade the luminosity analysis since with this prescaling the statistical error is still smaller than that for any other measured process.

4.1.2 Event selection

The main background in the luminosity event sample consists of beam-gas events. In order to distinguish between Bhabha events and this background all contiguous crystals with a measured energy greater than 250 MeV are grouped into clusters. The measured energy deposition pattern of all clusters is compared with the known shape for electromagnetic showers; overlapping clusters can thus be split up. The measured energy of all clusters is then corrected for lateral losses again using the known shower shape to obtain the best energy precision.

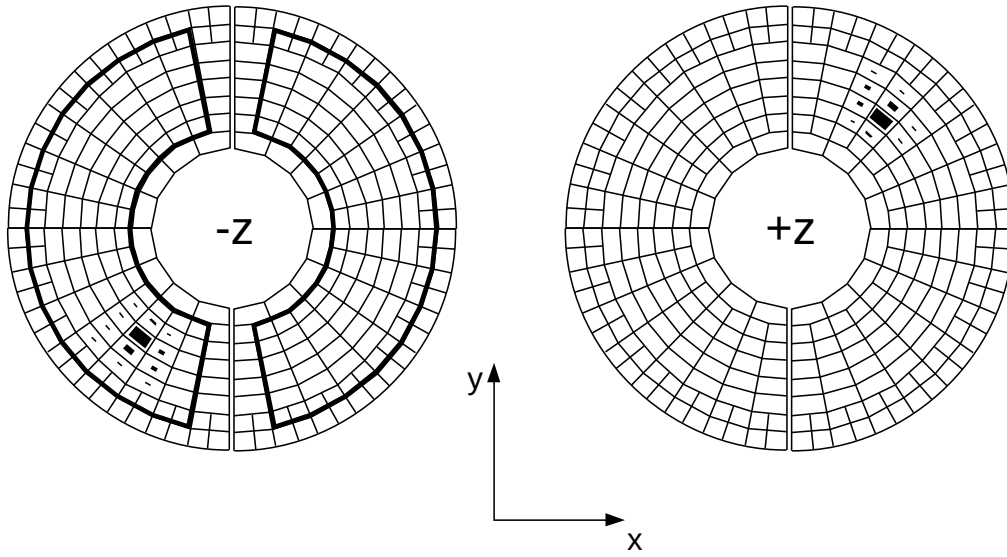


Figure 4.1: A typical Bhabha event as seen in the luminosity monitor. Shown is a schematic front view of the BGO crystals; the sizes of the black areas denote the energies deposited per crystal.

Subsequently all clusters in one and the same calorimeter are ordered by energy. A vectorial sum of all cluster momenta is made; this summation stops when the difference of the summed energy and the beam energy is minimal. The resulting superclusters (henceforth denoted as ‘clusters’) are used for further analysis. A typical ‘lumi’ event is shown in fig. 4.1.

To reduce the effects of misalignments [52] (detector offsets, beam offsets or beam tilts) the integrated luminosity is defined as the average of two samples with asymmetric fiducial volume cuts: one with tight cuts at the $+z$ side and no cuts on the $-z$ side, the other with tight cuts at the $-z$ side and no cuts on the $+z$ side. Calculations show that by using such an average, the effects of the experimentally measured offsets (which are within ± 2 mm) on the luminosity become negligible.

In more detail, the event selection criteria for each of the two samples are as follows:

1. *Fiducial volume*: on the tight fiducial volume side the reconstructed impact point must be at least one crystal away from the calorimeter edge, both in ϕ ($|\phi - 90^\circ| > 11.25^\circ$, $|\phi - 270^\circ| > 11.25^\circ$) and in θ ($30.92 < \theta < 64.41$ mrad for the 1990 data taking period and $29.56 < \theta < 61.66$ mrad for the 1991 data taking period.). There are no requirements on the impact point on the opposite side.
2. *Energy*: the energy on one side must exceed $0.8E_b$; the energy on the other side must be greater than $0.4E_b$.
3. *Coplanarity*: the event coplanarity angle $\Delta\phi$ must satisfy $|\Delta\phi - 180^\circ| < 10^\circ$.

The remainder of the background in the sample is estimated by interpolating the sideband values of the coplanarity distribution, $10^\circ < |\Delta\phi - 180^\circ| < 30^\circ$, into the acceptance region; this is done on a run-by-run basis.

4.1.3 Theoretical cross section

The theoretical Bhabha cross section corresponding to the fiducial volume covered by the luminosity monitor is determined using the BHLUMI Monte Carlo program [53]. This program features an $\mathcal{O}(\alpha)$ Yennie-Frautschi-Suura (YFS) [54] exponentiation to simulate the effects of multiple soft photon emission. The precision of the program was verified using a combination of two separate generators: one (OLDBIS) giving a full first order calculation, and a second (LUMLOG) calculating the $\mathcal{O}(\alpha^2 L^2)$ and $\mathcal{O}(\alpha^3 L^3)$ corrections. Here L indicates the so-called ‘large logarithms’ of the form $\ln(|t|/m_e^2)$ typical for bremsstrahlung in the t channel. An agreement to better than 0.25% was found.

Moreover, the cross section generated by the BHLUMI program was also compared with the result given by the BABAMC program [55]. This is an $\mathcal{O}(\alpha)$ Monte Carlo generator which uses an analytical calculation in order to estimate the $\mathcal{O}(\alpha^2 L^2)$ contributions [56]; it claims a theoretical uncertainty of 0.5%. The two results agree within 0.2% [57].

For the purpose of the L_3 luminosity measurement, Monte Carlo events are generated with a e^- scattering angle $20 \text{ mrad} < \theta < \pi - 20 \text{ mrad}$ and with a center-of-mass energy $\sqrt{s} = 91.18 \text{ GeV}$. For other, off-peak, energies the cross section is scaled using a $1/s$ dependence. In addition the energy dependence of the (small) electroweak interference effects is also accounted for.

The visible cross section (i.e. the cross section obtained after application of selection criteria) is determined, for the peak energy and for a nominal position of the luminosity monitor, using a sample of 800,000 fully simulated BHLUMI events. Subsequently this cross section is corrected at the generator level for the actual

energy and the position of the luminosity monitor using an approximation of the selection criteria. The visible cross sections corresponding to the luminosity monitor positions valid for the 1990 and 1991 data taking periods respectively, are found to be

$$\begin{aligned} & (84.7 \pm 0.1) \text{ nb}, \quad (1990) \\ & (90.3 \pm 0.1) \text{ nb}. \quad (1991) \end{aligned}$$

The uncertainties quoted are due to the finite size of the Monte Carlo sample only.

Because the analysis does not distinguish electrons and photons, there is an irreducible background from the process $e^+e^- \rightarrow \gamma\gamma(\gamma)$. Therefore the contribution from this process has to be added to the visible Bhabha cross section. This contribution is however very small and, like the small-angle Bhabha scattering itself, governed by QED; the increase of the visible cross section is calculated to be 0.02% only [58].

4.1.4 Systematic errors

The same sample of fully simulated BHLUMI events was used to estimate the stability of the measured luminosity against a variation of the selection criteria. The resulting variations are taken as a measure of the systematic uncertainties due to the event selection.

A summary of the systematic uncertainties on the luminosity determination is given in table 4.1. The uncertainties are the same for the 1990 and 1991 data taking periods; the theoretical uncertainties for both periods are assumed to be fully correlated.

Source of uncertainty	Contribution to $\Delta\mathcal{L}$ (%)
Detector geometry	0.4
Event selection criteria	0.3
Monte Carlo statistics	0.1
Theory	0.3
Total systematic uncertainty	0.6

Table 4.1: *Systematic uncertainties in the luminosity measurement.*

Analyses suffer from periods during which a degraded detector performance affects the systematic uncertainties. It is therefore desirable to be able to discard such periods from the analysis. The flexibility necessary to meet these requirements is provided by the so-called LUMLIST file. This file contains, on a run-by-run basis, the results of the luminosity analysis, together with the number of events accepted

as ‘luminosity events’. The latter serves to calculate the statistical error on the luminosity for any given data sample.

4.2 Systematic Uncertainties

The selection of hadronic events was described in the previous chapter. Also the efficiencies obtained for the relevant physics background sources were given there. Taking into account also the luminosity measurements described above, the cross sections are computed from these numbers using the formula

$$\sigma_{\text{had}} = \frac{1}{\epsilon_{\text{h}}}(N_{\text{h}}/\mathcal{L} - \epsilon_{\tau\tau}\sigma_{\tau\tau} - \epsilon_{\text{ee}}\sigma_{\text{ee}} - \epsilon_{\gamma\gamma}\sigma_{\gamma\gamma}) \quad (4.3)$$

where N_{h} and $\mathcal{L}$ are integrated over the complete data sample for each measurement point. This method allows a Gaussian treatment of the statistical uncertainties in these numbers. These uncertainties are simply calculated from the observed numbers of hadron and luminosity events, as

$$\Delta N_{\text{h}}/N_{\text{h}} = 1/\sqrt{N_{\text{h}}}, \quad \Delta \mathcal{L}/\mathcal{L} = 1/\sqrt{N_{\text{lumi}}}. \quad (4.4)$$

The selection criteria are devised to yield efficiencies that are as energy independent as possible. For this reason, only events generated at $\sqrt{s} = 91.25$ GeV are used in the determination of the efficiencies.

Besides the statistical errors, also the systematic ones need to be considered. For instance, the statistical uncertainties on the Monte Carlo efficiencies (discussed in section 3.2.6) are treated as systematic uncertainties on the cross sections: they are fully correlated because within one data taking period, always the same Monte Carlo sample is used. In the following, studies of other systematic effects are presented.

4.2.1 Beam energy spread

As mentioned in chapter 2, the spread of the particle energies within one bunch is approximately 50 MeV. One therefore does not measure the cross section at a single $\sqrt{s}$ point but the convolution of the cross section curve with a Gaussian resolution function. This introduces a small bias in the measured cross section which can be calculated if the theoretical shape of the cross section is known.

One proceeds in the following way: theoretical cross sections are calculated in a ‘model independent’ approach (as discussed in chapter 5). The convoluted cross section (which is obtained by numerical integration) is compared with the cross section in the absence of a beam energy spread. The ratio of these two results then gives the effect of the convolution which is to be corrected for. The correction is found to be largest at the peak of the resonance (+0.13%). Varying independently the theoretical input parameters within reasonable ranges one finds that the variation

of the correction is negligibly small ($\sim 0.01\%$) so that, to a good approximation, the correction factors can be applied *a priori*.

A similar effect is caused by the fact that there is a small spread in the fill-to-fill central beam energy values. However, this spread is significantly smaller than the aforementioned beam energy spread (the rms of the spread weighted by the fill integrated luminosity ranges from 10 to 20 MeV). One can show that roughly, the correction to be applied is proportional to the square of this rms energy spread; hence its effect is negligible in all cases.

4.2.2 Trigger efficiencies

In principle, the calculation of the selection efficiency as performed in the previous chapter does not necessarily lead to a correct determination of the *overall* efficiency. One factor which is not taken into account is the effect of inefficiencies in the L₃ trigger system. The trigger system is designed to be fully efficient, but problems with e.g. dead trigger channels or wrong trigger pedestals affect its performance. This could be taken into account via a Monte Carlo simulation of the trigger system, but this feature was unavailable for the MC programs used in the present analysis.

In order to achieve an as high as possible efficiency, the L₃ trigger system was designed such that for many classes of events more than one trigger is highly efficient. In the case of hadronic events, the main triggers are the energy trigger, the TEC trigger and the scintillator multiplicity trigger (see chapter 2). In addition to yielding a high combined trigger efficiency, this redundancy also allows an estimate of the combined efficiency from the selected data itself.

In order to arrive at such an efficiency estimate, it is necessary to know both the trigger efficiencies for the individual triggers and the way in which they are correlated. To a first approximation one can assume that there are no statistical correlations between the triggers, and that only geometrical effects resulting from the coverage of polar angular regions by more than one subdetector have to be accounted for. Thus a treatment of the efficiencies, bin-by-bin in the reconstructed event polar angle, is expected to yield an unbiased result. For each polar angle bin seven event categories are considered, following the combination of triggers by which the events are accepted. The eighth case (events accepted by none of the three above triggers) is not considered, since most of these events will not be accepted by the event selection criteria anyway. A log likelihood is defined:

$$\ln(L) = \sum_{i=1}^7 \ln P(\mu_i; n_i) \quad (4.5)$$

where $P(\mu_i; n_i)$ is the Poisson probability of finding n_i events in category i , when the expected number is μ_i ; μ_i depends on the efficiencies ϵ_E (energy trigger), ϵ_T (TEC trigger) and ϵ_S (scintillator multiplicity trigger). The quantity $\ln(L)$ is maximised as a function of ϵ_E , ϵ_T and ϵ_S . To allow for a possible incompleteness of the sample,

the total number of events per bin is also included as a fit parameter. The results, i.e. the trigger efficiencies weighted by the number of fitted events in each bin, are presented in table 4.2. Figure 4.2 shows the fitted inefficiencies as a function of the reconstructed $|\cos \theta_t|$ (the sign is irrelevant for this study).

	1990	1991
ϵ_E (%)	99.83 ± 0.02	99.31 ± 0.02
ϵ_T (%)	94.60 ± 0.09	92.63 ± 0.07
ϵ_S (%)	94.85 ± 0.09	92.75 ± 0.06
Combined	99.98 ± 0.01	99.93 ± 0.01

Table 4.2: Trigger efficiencies relevant for hadronic events. The numbers are obtained from fits to selected 1990 and 1991 data. The quoted errors are statistical only.

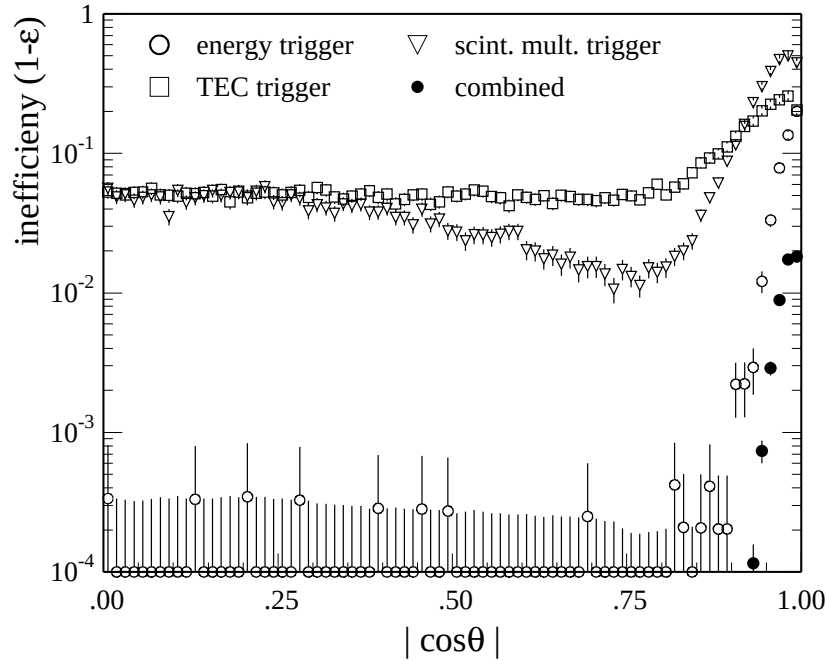


Figure 4.2: Individual and combined trigger inefficiencies as a function of the reconstructed event polar angle (1991 data).

The statistical independence of the various triggers can be tested, using the data sample itself. The efficiency of each trigger can be calculated from three independent samples: the events that were accepted by both of the other two triggers, or the

events accepted by just one of these. The resulting efficiencies are in reasonable agreement both with each other and with the values obtained from the fit, indicating that there are no large statistical correlations. This is in fact to be expected since each of the three mentioned triggers uses different subdetector information.

In any case, the region where the trigger is most inefficient (i.e. the region of small polar angles) is also the region where the data selection is most inefficient, due to the fact that the beam pipe does not allow detection of large parts of the event. Furthermore, the combined trigger efficiency is nearly completely determined by the energy trigger efficiency because this is by far the most efficient trigger. Given the fact that the energy trigger uses the same solid angle as the data selection procedure, and that the energy trigger thresholds are much lower than the minimum energy required by the data selection, it is very improbable for an event to be accepted by the selection procedure and not to fire the trigger system. For this reason, no correction for the trigger inefficiency is applied.

4.2.3 Fill-to-fill cross sections

In the search for systematic effects in the measured cross sections, one of the things that can be checked is the time dependence of these cross sections. One way to perform this check is to consider the cross sections obtained on a fill-by-fill basis. For each individual energy point, the cross section is calculated for each fill corresponding to this point; the expected number of events is calculated from the average cross section (as given in table 4.5) and the fill-integrated luminosity. The statistical error on the per-fill cross section is derived from this expected number of events rather than from the measured number of events in order to avoid introducing statistical biases. In addition a correction is applied for the differences in the fill-to-fill $\sqrt{s}$ values using the theoretical cross section. The result is shown (for the case of the 1991 data) in figure 4.3. The bottom part of this figure shows no significant variation as a function of time; the top part shows the fluctuations of the fills around the average cross section.

In case of purely statistical fluctuations, one expects the distribution describing the fluctuations around the average measured value to have an rms width of one. The calculated numbers for the 1990 and 1991 data taking periods are 0.99 ± 0.06 and 0.96 ± 0.06 respectively. These numbers being compatible with one, there is no reason to assume that time dependent effects play a significant role.

4.2.4 Non-resonant background

In the context of this thesis, the term ‘non-resonant background’ denotes all background processes which do not exhibit the Breit-Wigner resonance curve typical for Z exchange. This leaves the background from the two-photon process (discussed in section 3.2.2) and the one caused by ‘non-physics’ processes (mainly beam-gas

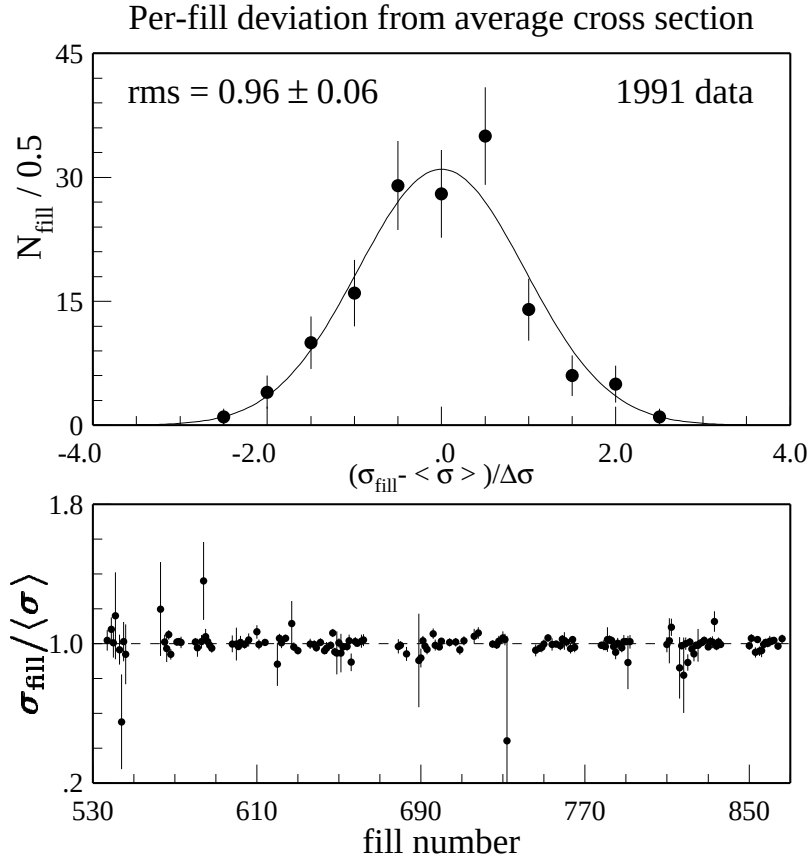


Figure 4.3: Fill-to-fill cross sections for the 1991 data taking period. Top part: deviations from the measured average cross sections. Bottom part: relative cross sections as a function of fill number.

events).

These two sources of background form a special problem for the present analysis. As stated in the previous chapter, the cross section of the two-photon process is not precisely known. The beam-gas background cannot even be estimated on the basis of a cross section. This implies that special care is needed in the treatment of these sources of background.

The situation is made even more difficult by the fact that using calorimetric information alone it is practically impossible to distinguish between two-photon events and beam-gas background. In principle it should be possible to make a distinction based on the fact that the beam-gas background is essentially a lepton-nucleon scattering process which can take place anywhere in the vacuum pipe, whereas the two-photon background is a genuine ‘physics’ background initiated by an e^+e^- interaction. This feature has been exploited in other experiments [59] to identify the beam-gas background using:

- the z coordinate of the reconstructed event vertex; one expects the two-photon z distribution to peak at the average vertex position, whereas the beam-gas distribution is expected to be flat;
- the dE/dx information; this is used to identify deuterons in the event, particles which are expected to be frequently produced in beam-gas events.

In this thesis, no attempt is made to actually separate the two-photon and beam-gas background contributions. Assuming that the E_{vis} and $E_{\parallel}$ distributions for these two are about the same (an assumption which cannot actually be verified) they are just discussed in terms of an *effective* two-photon cross section.

The simplest approach would be to fit a $\sqrt{s}$ independent cross section to this effective two-photon background. This would be a correct assumption if the beam-gas contribution can be disregarded. One would then proceed by performing a fit to the $E_{\parallel}$ distribution which, as confirmed by a visual inspection, gives a clean signal. An analysis along these lines shows however that the data are incompatible with this hypothesis: the effective cross section fitted to the 1990 data is significantly higher than the one fitted to the 1991 data; the assumption of a single, time independent cross section does not yield a good fit. This is consistent with the fact that the quality of the vacuum around the interaction points was better in the 1991 data taking period than in the 1990 one.

One way to taken into account the beam-gas contribution is to the drop the requirement that the effective two-photon cross section be $\sqrt{s}$ independent. One must then fit the background for each measurement point separately. The statistics on the off-peak points does not allow one to do this on the basis of the $E_{\parallel}$ distributions; thus the low-energy tails of the E_{vis} distributions are used for this purpose. The fit results for each $\sqrt{s}$ point are presented in table 4.3 (note that the actual $\sqrt{s}$ values are not the same for the two periods).

$\sqrt{s}$	$\sigma_{\gamma\gamma}$ (nb)	
	1990	1991
peak -3 GeV	3.8 ± 0.6	3.8 ± 0.2
peak -2 GeV	3.4 ± 0.6	3.8 ± 0.3
peak -1 GeV	4.2 ± 0.9	3.4 ± 0.4
peak	4.4 ± 0.9	2.7 ± 0.1
peak +1 GeV	3.1 ± 0.6	2.7 ± 0.5
peak +2 GeV	3.4 ± 0.6	3.5 ± 0.5
peak +3 GeV	3.8 ± 0.6	2.8 ± 0.3

Table 4.3: *Fitted effective two-photon cross sections at different energy points, for the 1990 and 1991 data.*

For all these results, the quoted uncertainties take into account both the statistical and systematic errors (the latter being estimated by varying the fit range under the condition that a good fit is maintained). As can be seen from table 4.3, for the 1990 and the 1991 data considered separately the fitted effective cross section varies only little with $\sqrt{s}$. This is to be expected from the fact that within the 1990 or 1991 data taking period, the quality of the vacuum inside the beam pipe near the L_3 interaction point, although different for the two periods, was relatively constant.

As an alternative method, instead of assuming all of the non-resonant background to look like a two-photon process, an exponential function is fitted to the low-energy tail of the E_{vis} distribution. Extrapolating the fitted function into the acceptance region then yields an estimate of the accepted background cross section. The results obtained are typically 30 pb higher than those using the two-photon method for the 1990 data, and typically 20 pb lower in the case of the 1991 data. Figure 4.4 shows for one energy point the two fits.

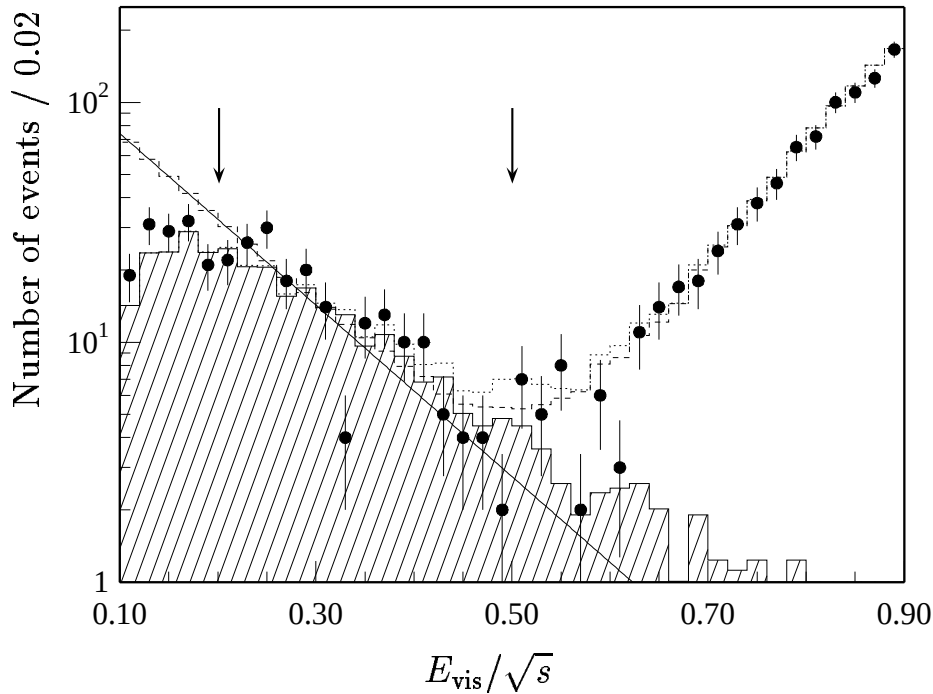


Figure 4.4: *Fit of the nonresonant cross section to the lowest-energy point for the 1991 data. The hatched area and the dotted line denote the nonresonant and total Monte Carlo predictions in the first method; the straight and dashed lines denote the fitted exponential function and total Monte Carlo predictions in the second fit method. The arrows denote the fit region.*

The fitted effective two-photon cross sections of table 4.3 are used to correct the measured hadronic cross sections. The differences with the second method (i.e. 30

pb for the 1990 data and 20 pb for the 1991 data) are taken as a measure of the systematic error due to this background subtraction. It is assumed that there exists no correlation between the uncertainties for different data points.

4.2.5 Selection criteria

In section 3.2.4, the reasons for the choice of selection criteria were given. However, these reasons only yield the approximate cut values; this imprecision is of minor importance if the measured cross sections do not change significantly when varying the cuts within a reasonable range. This can be tested by varying any of the cuts while keeping the others fixed.

The difference in cross section obtained for two different sets of cuts can be approximated by

$$\delta\sigma/\sigma \approx \delta N_{\text{ev}}/N_{\text{ev}} - \delta\epsilon_{\text{h}}/\epsilon_{\text{h}} - \frac{1}{\sigma\epsilon_{\text{h}}}(\sigma_{\text{ee}}\delta\epsilon_{\text{ee}} + \sigma_{\tau\tau}\delta\epsilon_{\tau\tau} + \sigma_{\gamma\gamma}\delta\epsilon_{\gamma\gamma}) \quad (4.6)$$

where δ denotes the difference caused by going from the looser set of cuts to the tighter one. This approximation is good if the background contamination is small and the signal efficiency high. In addition the variation in the sample induced by the cut change should also be small since eq. 4.6 neglects several higher order terms. From this formula, an approximate expression for the uncertainty on $\delta\sigma$ can be derived:

$$\begin{aligned} \Delta^2(\delta\sigma/\sigma) &\approx \Delta^2(\delta N_{\text{ev}}/N_{\text{ev}}) + \Delta^2(\delta\epsilon_{\text{h}}/\epsilon_{\text{h}}) \\ &+ \frac{1}{(\sigma\epsilon_{\text{h}})^2} \left[\Delta^2(\delta\epsilon_{\text{ee}})\sigma_{\text{ee}}^2 + \Delta^2(\delta\epsilon_{\tau\tau})\sigma_{\tau\tau}^2 + \Delta^2(\delta\epsilon_{\gamma\gamma})\sigma_{\gamma\gamma}^2 \right]. \end{aligned} \quad (4.7)$$

where Δx denotes the uncertainty in quantity x . Binomial errors are used to estimate the various sources of uncertainties. Figure 4.5 shows, for the 1991 peak point, the relative variations in the cross section, as a function of the most important cut variables.

In the ranges considered, the observed differences with the central cross section point are in general smaller than 0.2% for the 1990 data and smaller than 0.1% for the 1991 data. Conservatively, a systematic error due to uncertainties in the selection is assigned of 0.1% for the 1990 data and 0.05% for the 1991 data; no correlation is assumed between the errors for different data points.

4.2.6 Random background

As mentioned in section 3.2.5, a correction is applied to estimate the effects of additional noise clusters due to random background. This is done using the distribution of the number of clusters observed in beam-gate events. This method in principle corrects for any possible effect. However, it also entails a certain bias because these

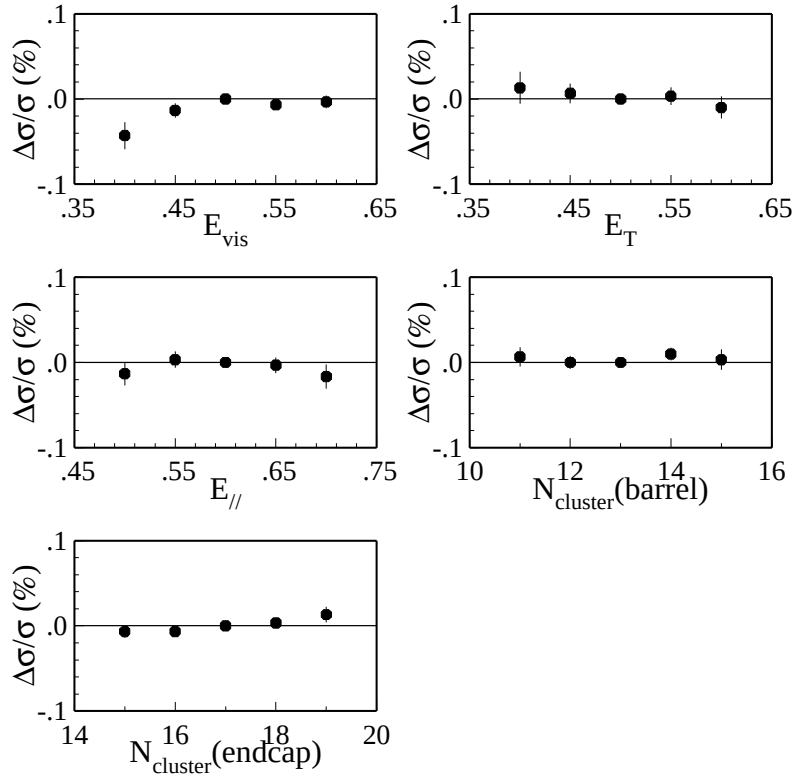


Figure 4.5: Relative variations of the measured cross section for the peak point (1991 data).

beam-gate events contain uranium noise. Convoluting the beam-gate cluster distribution with the N_{cluster} distribution found in Monte Carlo data one thus introduces some double counting. In order to estimate the possible systematic effect this might have on the measured cross section, the following procedure is followed:

- Once more, the N_{cluster} distribution for beam-gate events is made; however, this time any cluster containing an HCAL cluster constituent is disregarded. It is assumed to be very improbable (since the ECAL is shielded against uranium noise) for uranium noise to induce ECAL bumps without creating HCAL clusters as well. One thus obtains a lower limit on the number of clusters induced by sources other than uranium noise. The distributions involved for the 1991 data are shown in figure 4.6.
- The inclusion, respectively exclusion, of the beam-gate clusters containing an HCAL constituent reflects itself in a change of the Monte Carlo efficiencies. Since *a priori* neither of the two methods is expected to give completely correct results, the approach taken is to average the two efficiencies and to consider

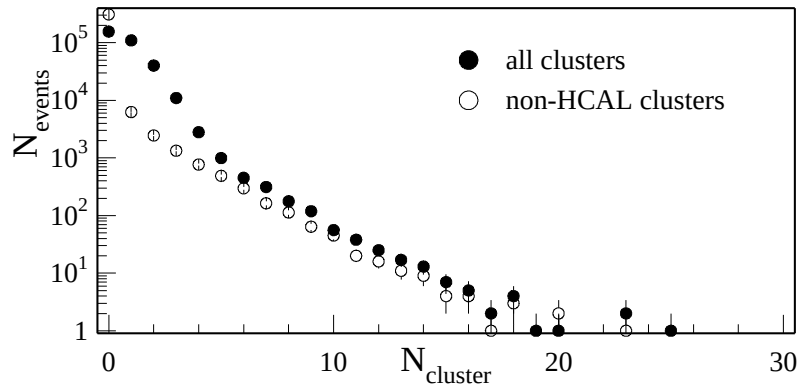


Figure 4.6: N_{cluster} distribution in beam-gate events (1991 data).

half their difference as a systematic uncertainty. The half differences obtained for the 1990 and 1991 data are shown in table 4.4. As can be seen from this table, the effect on $\epsilon_{\tau\tau}$ is the most pronounced. Since these systematic errors all follow from the same uncertainty in the use of beam-gate clusters, they are fully correlated.

	1990	1991
$\Delta\epsilon_h$	0.034%	0.028%
$\Delta\epsilon_{\tau\tau}$	1.828%	1.002%
$\Delta\epsilon_{ee}$	0.318%	0.003%
$\Delta\epsilon_{\gamma\gamma}$	0.010%	0.003%

Table 4.4: Systematic uncertainties on the various efficiencies, for the 1990 and 1991 data. It should be stressed that these errors are fully correlated.

The translation of the uncertainties of table 4.4 into uncertainties on the hadronic cross sections can be simplified by the following two approximations. First, since the $e^+e^- \rightarrow \tau\tau(\gamma)$ process is an s -channel process, the $\tau^+\tau^-$ cross section exhibits to a good approximation the same resonance behaviour as the $q\bar{q}$ cross section. For the present purpose it therefore assumed to scale with the $q\bar{q}$ cross section as $\sigma_{\tau\tau} \approx 0.048\sigma_{\text{had}}$; the proportionality factor of 0.048 represents the ratio of the measured peak cross sections [29]. Second, the energy dependence of the Bhabha cross section and the (effective) two-photon cross section is neglected. These approximations

allow to summarise the consequences of the efficiency errors for the total systematic error on σ_{had} in terms of two contributions:

1. a *constant* random background error contribution due to the combination of $\Delta\epsilon_{ee}$ and $\Delta\epsilon_{\gamma\gamma}$; this yields 43 pb for the 1990 and 1 pb for the 1991 data;
2. a random background contribution which *scales* with σ_{had} and is due to the combination of $\Delta\epsilon_h$ and $\Delta\epsilon_{\tau\tau}$; for 1990 and 1991 one finds the value of 0.13% and 0.08% respectively.

4.2.7 Fragmentation

The JETSET program provides, as shown above, a good description of the data. However, it contains many parameters that were in fact tuned to reproduce these data. It is therefore desirable to have an independent check of the selection efficiency obtained.

For this purpose, the HERWIG program [60] is used. This program features cluster fragmentation in conjunction with a Parton Shower algorithm (see section 3.2.1). In the branching process it also takes into account interference effects. Initial-state photon radiation is not implemented, but this is not expected to have a big effect since most of the initial-state radiation is soft.

A sample of 290,000 events was generated at $\sqrt{s} = 91.25$ GeV, corresponding to the 1991 detector geometry. One finds that in general, the HERWIG program describes the data less well than JETSET, as illustrated in figure 4.7. The efficiency obtained using this program is 0.15% lower than that obtained using the JETSET program; this number is used as an additional (correlated) systematic error following from the specific description used for the fragmentation process. Lacking a sizeable event sample corresponding to the 1990 detector geometry, the same number is assumed for the 1990 data.

4.3 Results

Taking into account the previously described systematic studies, the cross sections and corresponding statistical and systematic uncertainties are given in table 4.5. It should be noted that for the 1991 data taking period there are two data points for $\sqrt{s} \approx M_Z$, corresponding to the data taken before and after the first depolarisation calibration.

Making the same approximations as in section 4.2.6, the background in the selected event samples can be summarised as follows:

- The process $e^+e^- \rightarrow e^+e^-(\gamma)$ causes a $\sqrt{s}$ independent background of (46 ± 43) pb in the 1990 data and (1 ± 1) pb in the 1991 data;

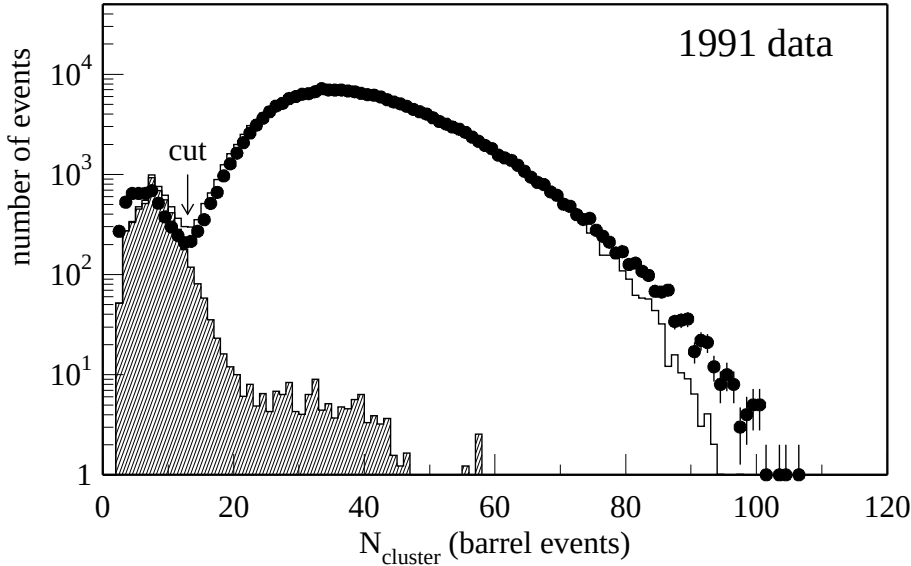


Figure 4.7: Comparison between 1991 data and corresponding Monte Carlo predictions, of the observed cluster multiplicity for endcap events. The hatched area represents the contribution expected from background processes; the second line inside this area denotes the same but with the two-photon contribution omitted. The simulated $e^+e^- \rightarrow \text{hadrons}$ events were generated using the HERWIG program.

- The non-resonant processes induce a $\sqrt{s}$ independent background of (40 ± 2) pb, in both the 1990 and 1991 data;
- The process $e^+e^- \rightarrow \tau\tau(\gamma)$ leads to a background contamination of $(0.34 \pm 0.09)\%$ in the 1990 data and $(0.19 \pm 0.05)\%$ in the 1991 data.

The uncertainties on these numbers originate from the finite sizes of the Monte Carlo samples as well as from the uncertainties in the treatment of random background.

Table 4.6a presents a break-down of the systematic uncertainties obtained for the 1990 and 1991 data samples. The statistical uncertainties in the efficiency estimates are translated into systematic uncertainties on the hadronic cross sections using the same simplification as in section 4.2.6.

Taking together all effects, one thus again distinguishes two error contributions: one which is constant as a function of $\sqrt{s}$ and one which scales with the measured hadronic cross section. Accounting in addition for the fact that these errors are partially correlated for the various energy points, one obtains the ‘total’ systematic errors of table 4.6b. To obtain the complete systematic uncertainty for a given $\sqrt{s}$ point, one then makes the appropriate linear and quadratic combinations depending on the presence or absence of correlations. It should however be stressed that

in combining data points, there are always additional correlations which should be taken into account. To a good approximation, it can be assumed that only the correlated scale error is also correlated for the 1990 and 1991 data; the other categories are uncorrelated.

1990 data

$\sqrt{s}$ (GeV)	N_{events}	$\mathcal{L}$ (nb $^{-1}$)	σ_{had} (nb)
88.232	1525	339.36	$4.453 \pm 0.118 \pm 0.058$
89.236	3600	422.87	$8.516 \pm 0.149 \pm 0.064$
90.237	6147	330.93	$18.680 \pm 0.261 \pm 0.081$
91.230	79157	2623.98	$30.445 \pm 0.125 \pm 0.104$
92.226	8182	374.36	$22.010 \pm 0.272 \pm 0.087$
93.228	5922	480.30	$12.378 \pm 0.172 \pm 0.070$
94.223	3752	465.69	$8.056 \pm 0.138 \pm 0.063$
Total	108285	5037.49	

1991 data

$\sqrt{s}$ (GeV)	N_{events}	$\mathcal{L}$ (nb $^{-1}$)	σ_{had} (nb)
91.253	155084	5124.09	$30.444 \pm 0.095 \pm 0.058$
88.480	4050	781.89	$5.171 \pm 0.086 \pm 0.022$
89.470	8528	846.74	$10.091 \pm 0.120 \pm 0.027$
90.228	14334	793.29	$18.152 \pm 0.178 \pm 0.038$
91.222	90621	3010.79	$30.325 \pm 0.129 \pm 0.058$
91.967	16058	657.61	$24.567 \pm 0.240 \pm 0.048$
92.966	10863	758.24	$14.386 \pm 0.159 \pm 0.033$
93.717	7147	706.15	$10.157 \pm 0.133 \pm 0.027$
Total	306688	12678.80	

Table 4.5: Measured cross sections for the 1990 and 1991 data. The first errors represent the statistical uncertainties on the luminosity and on the number of selected hadronic events; the second ones represent the systematic uncertainties. It should be noted that these systematic uncertainties are correlated for the various data points; in addition, the 0.6% systematic uncertainty on the luminosity measurement is not taken into account.

A last remark is that the 0.6% systematic uncertainty on the luminosity measurement is an uncertainty which is common to all cross section measurements; this error is treated separately in the fits performed in the next chapter and as such is not included in table 4.6.

a: Source of uncertainty	1990 data	1991 data
MC $q\bar{q}$ statistics	0.03%	0.01%
MC $\tau^+\tau^-$ statistics	0.01%	-
MC e^+e^- statistics	5pb	1pb
MC two-photon statistics	3pb	2pb
Non-res. background	30pb	20pb
Variation of cuts	0.10%	0.05%
Fragmentation	0.15%	0.15%
Random background (scaling)	0.13%	0.08%
Random background (constant)	43pb	1pb

b: Total systematic errors		
Correlated error (scaling)	0.20%	0.17%
Uncorrelated error (scaling)	0.10%	0.05%
Correlated error (constant)	43pb	2pb
Uncorrelated error (constant)	30pb	20pb

Table 4.6: *Systematic uncertainties for the 1990 and 1991 data. (a) individual contributions. (b) combination.*

Chapter 5

Analysis of Lineshape Data

The hadronic cross sections obtained in the previous chapter are the basis for the determination of various parameters. It is these measurements which dominate the precision of the two most ‘elementary’ parameters of the Z boson, its mass M_Z and its total decay width Γ_Z . However, there are other important Standard Model parameters than those just mentioned. In order to obtain more detailed information about the Standard Model parameters the results from other measurements have to be invoked. This chapter therefore also makes use of the results obtained for the leptonic cross sections and the forward-backward asymmetry measurements (see chapter 1). In principle one could also have included the results obtained for the τ polarisation, the $b\bar{b}$ cross section and the $b\bar{b}$ forward-backward asymmetry. For the present purpose the extra information has however been limited to the aforementioned.

Two approaches can be taken in interpreting the lineshape measurements. The first of these is a ‘model independent’ approach which aims at the determination of *effective* parameters (e.g. the partial and total Z decay widths); the results obtained using this method are (nearly) independent of the electroweak structure of the Standard Model. They can be used to test the compatibility of the measured data with several models, among which the Standard Model itself.

In the second approach one assumes that the Standard Model is strictly valid. The constraints imposed by this assumption can then be used to extract information about the parameters which are presently unknown within the context of the Standard Model.

The program used to compare the experimental data to the theoretical predictions is the ZFITTER [8] package. This program is based on a semi-analytical treatment of the various fermion pair production processes. It allows for the two approaches described above. The Born-like cross sections, and in the second approach, the electroweak and QCD corrections, are calculated analytically. In the

latter case, the complete $\mathcal{O}(\alpha)$ m_t -dependent terms as well as the leading $\mathcal{O}(\alpha\alpha_s m_t^2)$ and $\mathcal{O}(\alpha^2 m_t^2)$ terms are taken into account. Photonic corrections up to $\mathcal{O}(\alpha)$ and soft photon exponentiation are implemented using a structure function approach (see section 1.3.2). The photon phase space is restricted through kinematical cuts; the integration over the allowed photon phase space is performed numerically. The advantage of the applied structure function formalism is that it allows (in principle) to convolute known QED corrections with any underlying electroweak structure.

5.1 Model Independent Fits

As stated in section 1.3, the main effect of electroweak corrections is to replace the coupling constants g_V^f, g_A^f with *effective* coupling constants $\bar{g}_V^f, \bar{g}_A^f$. Due to these corrections, and to the QCD corrections in the case of hadronic final states, the most important effect (as far as the total cross section is concerned) is to modify the partial decay widths Γ_f (and hence the total Z decay width Γ_Z). Furthermore, the s -dependent width 1.27 comes into play.

For this reason, one can to a good approximation combine equations 1.13, 1.15 and 1.27 to obtain the total cross section for $f\bar{f}$ final states ($f \neq e$) in terms of the Z mass and its partial and total decay widths. The result is

$$\begin{aligned} \sigma^f(s) \approx & 12\pi \frac{\Gamma_e \Gamma_f}{M_Z^2} \frac{s}{(s - M_Z^2)^2 + s^2 \Gamma_Z^2 / M_Z^2} \\ & - \frac{2\sqrt{2}\alpha(M_Z^2)N_C^f}{3} G_\mu g_V^e g_V^f Q_f (1 + 2\mu_f) \frac{(s - M_Z^2)M_Z^2}{(s - M_Z^2)^2 + s^2 \Gamma_Z^2 / M_Z^2} \\ & + \frac{4\pi Q_f^2 \alpha^2 (M_Z^2) N_C^f}{3s} \end{aligned} \quad (5.1)$$

where it is assumed that QED and QCD corrections to Γ_e and Γ_f have already been absorbed into the definition of these quantities. Since no distinction is made between the various quark flavours, one adds the quark partial decay widths to obtain the total hadronic decay width Γ_{had} .

Evidently, the interference term in this formula cannot easily be written in terms of ‘directly observable’ quantities – in principle it could be written as a function of the positive and negative helicity decay widths $\Gamma_f^\pm$ [7], but the measurement of those helicity widths requires the determination of the spin of the fermions involved and was therefore not attempted. However, on the Z resonance this interference term is small. The approach taken in the ZFITTER program is to infer this term from a Standard Model calculation. For the purpose of this calculation only, top mass and Higgs mass values $m_t = 150$ GeV and $M_H = 250$ GeV are assumed. In principle this procedure introduces a small model dependence. Varying this term between zero and twice its nominal value changes M_Z by only a few MeV, and does not affect the widths at all [37].

From eq. 5.1 it is clear that also Γ_e is needed to obtain a value for Γ_{had} . There are thus two possibilities: one can either use the Standard Model prediction for Γ_e or use the measurement of this parameter in large-angle Bhabha scattering. To allow for a maximum sensitivity to Standard Model parameters the latter method is preferred. However, it should be stressed that also this approach entails some model dependence. As explained in section 1.2, the Bhabha cross section receives contributions from both s -channel and t -channel diagrams. In order to derive results pertaining to Z exchange, the t -channel and $(s-t)$ interference contributions are inferred from the Standard Model using the ALIBABA [61] program, and subtracted from the measured cross sections.

Since all leptons are much lighter than the Z boson, mass effects are expected to play only a negligible role in the total leptonic cross sections. If the coupling of the Z boson to all leptons is the same (which is the case in the Standard Model) one expects the same partial decay widths for all lepton flavours, i.e. $\Gamma_e = \Gamma_\mu = \Gamma_\tau$. To test this expectation a fit is performed in which all these partial widths are left as free parameters. A positive outcome of this test then gives credence to the hypothesis of a single leptonic width Γ_l for all the leptonic final states. One thus performs either a six parameter fit (without the assumption of ‘lepton universality’) using M_Z , Γ_Z , Γ_{had} and the various leptonic decay widths, or a four parameter fit using M_Z , Γ_Z , Γ_{had} and a universal leptonic partial decay width Γ_l .

The fits are performed employing a χ^2 method. The χ^2 is constructed using a covariance matrix containing the statistical and systematic uncertainties. This matrix allows one to properly take into account the correlations between the various measurements. The LEP $\sqrt{s}$ uncertainties enter through a Gaussian error propagation using the various theoretical cross sections as they depend on the parameters being fitted. The minimisation is performed using the MINUIT [62] package.

The results of the fits are given in table 5.1. A good fit is obtained for both hypotheses. The leptonic partial widths obtained in the six parameter fit are indeed compatible with each other, thus supporting the assumption of lepton universality. Furthermore, the fitted partial widths are compatible with the Standard Model predictions. The fit is also repeated omitting the LEP $\sqrt{s}$ uncertainties; one thus obtains an estimate of the relative importance of the experimental errors versus the LEP energy uncertainties. The results found are

$\begin{aligned} M_Z &= (91197 \pm 6 \pm 7) \text{ MeV}, \\ \Gamma_Z &= (2492 \pm 10 \pm 5) \text{ MeV} \end{aligned}$
--

where the first error is experimental and the second one due to the LEP energy uncertainties. From this result one can conclude that the measurement of M_Z would benefit from an improved LEP energy calibration.

As a cross-check, a fit to the hadronic cross sections is performed inferring the electronic partial width from the Standard Model, $\Gamma_e = 83.8 \text{ MeV}$. One obtains

Parameter	Fit method		Standard Model
	6 parameters	4 parameters	
M_Z (MeV)	91197 ± 10	91197 ± 10	
Γ_Z (MeV)	2492 ± 11	2492 ± 11	2489
Γ_{had} (MeV)	1753 ± 14	1747 ± 11	1737
Γ_e (MeV)	83.1 ± 0.7		83.8
Γ_μ (MeV)	83.4 ± 1.1		83.8
Γ_τ (MeV)	84.7 ± 1.3		83.6
Γ_l (MeV)		83.4 ± 0.5	83.8
χ^2/N_{DF}	51.3/54	52.8/56	

Table 5.1: Results of the fits to the Z mass and its partial and total decay widths. Six parameter fit: fit to hadronic and leptonic data without the assumption of lepton universality. Four parameter fit: fit to hadronic and leptonic data assuming lepton universality. The Standard Model predictions do not take into account variations of m_i and M_H .

again a good fit ($\chi^2/N_{\text{DF}} = 8.24/12$) with the results

$$\begin{aligned}
 M_Z &= (91196 \pm 12) \text{ MeV}, \\
 \Gamma_Z &= (2492 \pm 10) \text{ MeV}, \\
 \Gamma_{\text{had}} &= (1736 \pm 18) \text{ MeV}
 \end{aligned}$$

which (for M_Z and Γ_Z) is in good agreement with the previous fits. Evidently, the uncertainties on M_Z and Γ_Z are essentially determined by those on the hadronic cross sections. The fact that Γ_{had} is found to be somewhat lower is because the value of Γ_e obtained in the six parameter fit is lower than its Standard Model prediction. Figure 5.1 shows the four parameter fit of the hadronic line shape together with the experimental data. Because the statistical uncertainties on the hadronic cross sections are too small to be displayed on the lineshape plot itself, a blowup of the ratio $\sigma_{\text{meas}}/\sigma_{\text{fit}}$ is shown as well.

The tests of the Standard Model can be extended by checking its predictions up to the level of the individual effective coupling constants $\bar{g}_V^f, \bar{g}_A^f$. Note that, neglecting fermion masses, the total cross section only depends on the quadratic sum $(\bar{g}_V^f)^2 + (\bar{g}_A^f)^2$. Equation 1.22 indicates that in the forward-backward asymmetry case it is the product $\bar{g}_V^f \bar{g}_A^f$ which is measured. Thus, a combination of cross section and asymmetry measurements can be used to determine the magnitude of $\bar{g}_V^f, \bar{g}_A^f$. The results, given in table 5.2, again indicate lepton universality and compatibility

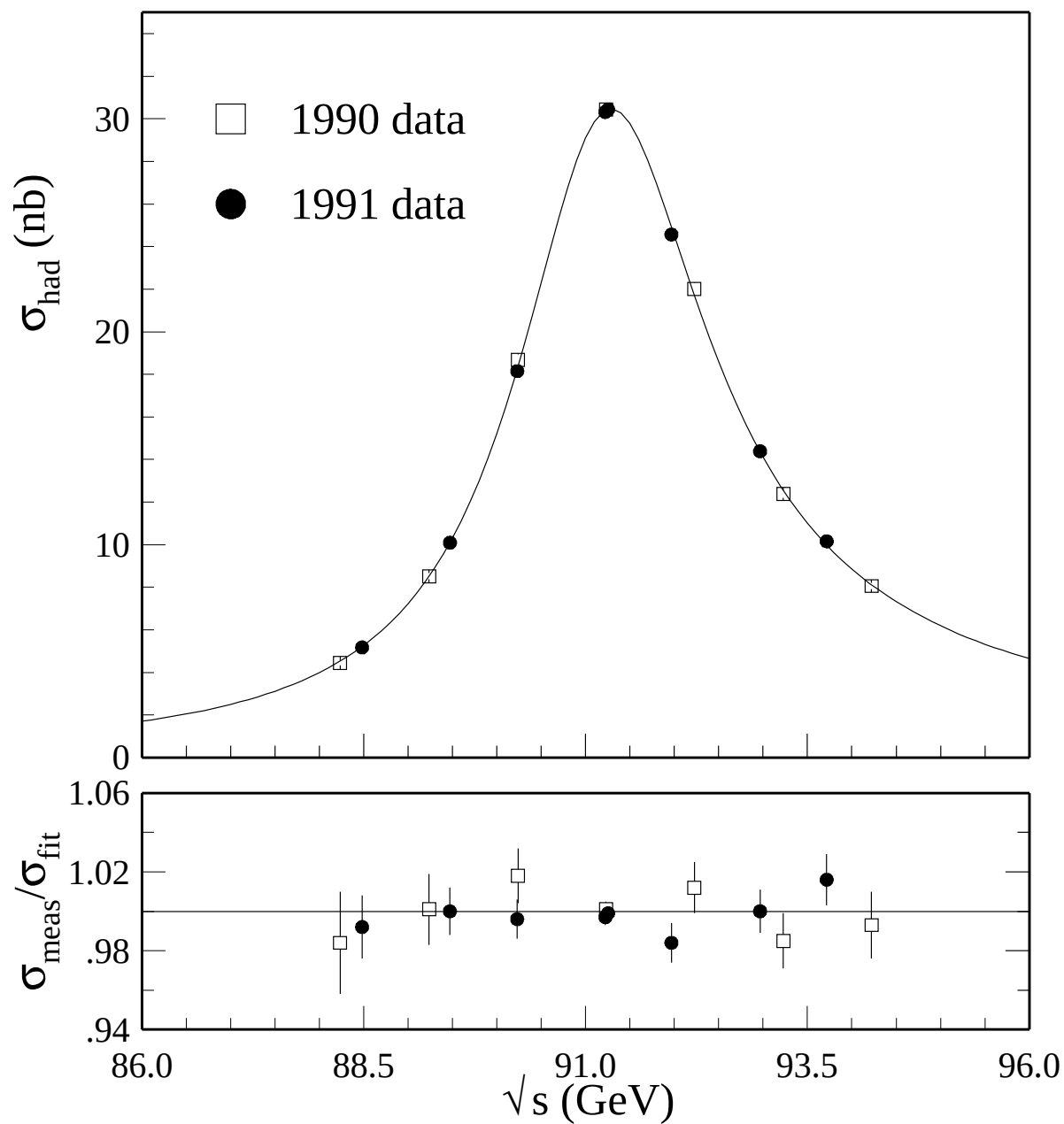


Figure 5.1: Comparison between the measured hadronic cross sections and the result of the four parameter fit. The error bars refer to the statistical uncertainties only.

with the Standard Model predictions. Note that the signs of the effective coupling constants $\bar{g}_V^f$, $\bar{g}_A^f$ are those predicted by the Standard Model; they have however been confirmed by neutrino-lepton scattering experiments (see e.g. the discussion in reference 3). Figure 5.2 compares the measured values of the coupling constants

Parameter	Fit method		Standard Model
	9 parameters	5 parameters	
M_Z (MeV)	91197 ± 7	91197 ± 10	
Γ_Z (MeV)	2492 ± 10	2492 ± 11	2489
$\bar{g}_V^e$	-0.0353 ± 0.0153		-0.0345
$\bar{g}_A^e$	-0.4987 ± 0.0021		-0.5008
$\bar{g}_V^\mu$	-0.0544 ± 0.0267		-0.0345
$\bar{g}_A^\mu$	-0.4978 ± 0.0041		-0.5008
$\bar{g}_V^\tau$	-0.0664 ± 0.0368		-0.0345
$\bar{g}_A^\tau$	-0.5005 ± 0.0059		-0.5008
$\bar{g}_V^l$		-0.0430 ± 0.0069	-0.0345
$\bar{g}_A^l$		-0.4990 ± 0.0016	-0.5008
χ^2/N_{DF}	82.4/96	84.3/100	

Table 5.2: *Effective weak coupling constants. Nine parameter fit: without assumption of lepton universality; five parameter fit: with lepton universality. The Standard Model predictions do not take into account variations of m_t and M_H .*

with the corresponding Standard Model predictions.

5.2 Number of Light Neutrino Species N_ν

The hadronic and leptonic final states constitute the channels which are directly measurable at LEP. Assuming that besides these final states, only light neutrinos (i.e. neutrinos with mass smaller than $M_Z/2$) with Standard Model couplings are produced, the widths obtained above can be used to derive the number N_ν of such neutrino species. In principle, one obtains N_ν by dividing the ‘invisible width’

$$\Gamma_{\text{inv}} = \Gamma_Z - \Gamma_{\text{had}} - 3\Gamma_l \quad (5.2)$$

by the Standard Model prediction for the neutrino width Γ_ν^{SM} . However, the systematic error on the luminosity directly enters in the result because both Γ_{had} and Γ_l depend on it. Moreover, Γ_ν^{SM} is sensitive to the unknowns of the Standard Model, m_t and M_H . Assuming central values $m_t = 150$ GeV and $M_H = 250$ GeV, and varying

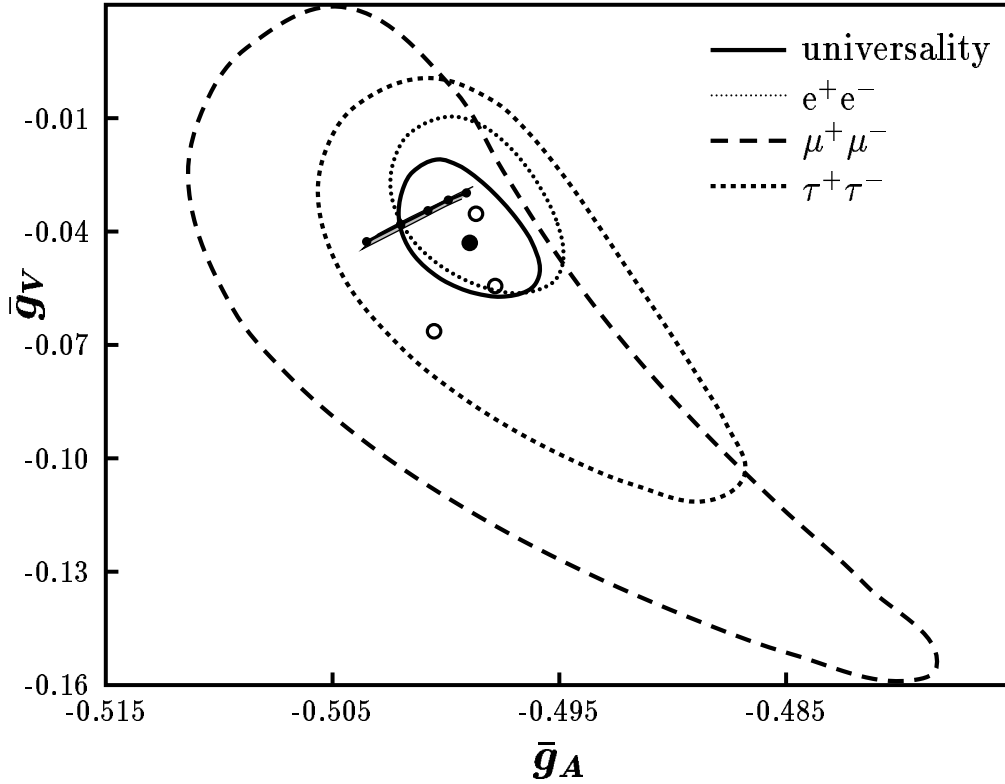


Figure 5.2: 95% C.L. contours in the $\bar{g}_V$ - $\bar{g}_A$ plane, either collectively assuming lepton universality (solid contour), or for each lepton species separately (dashed and dotted contours). The straight line denotes the Standard Model prediction for top masses varying between 50 GeV (the rightmost point) and 250 GeV (the leftmost point). The variation of this prediction for Higgs masses between 60 GeV and 1 TeV is denoted by the small band.

these two parameters between 50 and 250 GeV and 60 and 1000 GeV respectively, one obtains $\Gamma_\nu^{\text{SM}} = (167.0^{+1.9}_{-1.4})\text{MeV}$, i.e. variations of the order of 1%.

Fortunately, the ratio $(\Gamma_l/\Gamma_\nu)^{\text{SM}}$ is quite insensitive to these variations: varying m_t and M_H over the same ranges as before one finds $(\Gamma_l/\Gamma_\nu)^{\text{SM}} = 0.5015^{+0.0015}_{-0.0008}$, i.e. variations of 0.2% only. The number of light neutrino species is therefore better derived from

$$\begin{aligned}
 N_\nu &= \frac{\Gamma_{\text{inv}}}{\Gamma_l} \left(\frac{\Gamma_l}{\Gamma_\nu} \right)^{\text{SM}} \\
 &= \left(\frac{\Gamma_Z - \Gamma_{\text{had}}}{\Gamma_l} - 3 \right) \left(\frac{\Gamma_l}{\Gamma_\nu} \right)^{\text{SM}}.
 \end{aligned} \tag{5.3}$$

This leads to the result

$$\boxed{N_\nu = 2.98 \pm 0.06.} \tag{5.4}$$

The result is clearly compatible with the assumption that there exist three ‘light neutrino’ species with Standard Model couplings.

The ‘invisible width’ method can be compared with the direct ‘single photon’ method of measuring N_ν . This method entails the identification of the radiative process $e^+e^- \rightarrow \nu\bar{\nu}\gamma$ through the observation of single photons in the detector, and a comparison of the observed cross section with the Standard Model prediction leaving N_ν as a free parameter. The result of this analysis is [29]

$$N_\nu = 3.14 \pm 0.24(\text{stat.}) \pm 0.08(\text{syst.}) \quad (5.5)$$

in good agreement with equation 5.4, albeit with much larger uncertainties.

5.3 Strong Coupling Constant $\alpha_s(M_Z^2)$

As mentioned in section 1.3, one important difference between the hadronic and leptonic channels is that QCD corrections enter in the former, whereas they are absent in the latter. Thus the ratio $\Gamma_{\text{had}}/\Gamma_l$ carries information about these QCD corrections. Assuming Standard Model electroweak couplings, this ratio can be expressed as [63]

$$\begin{aligned} R_{\text{had}} &\equiv \Gamma_{\text{had}}/\Gamma_l \\ &= R_{\text{had}}^{\text{SM},0} \left(1 + 1.05 \frac{\alpha_s(M_Z^2)}{\pi} + (0.9 \pm 0.1) \left(\frac{\alpha_s(M_Z^2)}{\pi} \right)^2 - 13 \left(\frac{\alpha_s(M_Z^2)}{\pi} \right)^3 + \dots \right) \end{aligned} \quad (5.6)$$

where $R_{\text{had}}^{\text{SM},0}$ is the Standard Model prediction for R_{had} in the absence of strong corrections. Using the same variations of the input parameters m_t and M_H as before, this prediction is found to be $R_{\text{had}}^{\text{SM},0} = 19.96 \pm 0.03$. Using in addition the observed ratio

$$R_{\text{had}} = 20.96 \pm 0.15$$

equation 5.6 can be solved for $\alpha_s(M_Z^2)$ to obtain

$$\boxed{\alpha_s(M_Z^2) = 0.148_{-0.024}^{+0.012}}, \quad (5.7)$$

a value which is in reasonable agreement with those obtained from other α_s dependent quantities such as the event shape variables¹, $\alpha_s(M_Z^2)^{\text{shape}} = 0.125 \pm 0.003(\text{stat.}) \pm 0.008(\text{theor.})$, or the hadronic branching ratio B_τ of the τ lepton, $\alpha_s(M_Z^2)^{B_\tau} = 0.121 \pm 0.006(\text{expt.}) \pm 0.003(\text{theor.})$ [29].

¹Variables used are thrust, energy-energy correlation asymmetry, and the so-called heavy jet mass

5.4 Standard Model Fits

Having verified the agreement between observed quantities and the corresponding Standard Model predictions, one can now assume the validity of this model and try to estimate its parameters.

There are four input parameters in the Standard Model which are not (yet) very well established by other measurements: these are M_Z , m_t , M_H and $\alpha_s(M_Z^2)$.

Because the dependence of the measured quantities on M_H is very small (it is only logarithmic) no attempt is made to actually measure this parameter; instead it is fixed in the fit. Subsequently the effect of its variation is determined and considered as an additional (theoretical) uncertainty. In the following fits $M_H = 250$ GeV is used as a central value. As a lower limit 60 GeV is used, a value which roughly follows from the direct search for the Higgs boson in various channels [29]:

$$M_H > 52 \text{ GeV at } 95\% \text{ C.L.}$$

An upper limit of 1 TeV is imposed in order to prevent violations of unitarity at tree level in the scattering of longitudinally polarised gauge bosons [64], a requirement for any perturbative treatment in which the Higgs boson is involved to be valid.

A first fit is performed leaving M_Z , m_t and $\alpha_s(M_Z^2)$ as free parameters. In this case one obtains a good fit ($\chi^2/N_{\text{DF}} = 87/102$) with the results

$$\begin{aligned} M_Z &= (91196 \pm 9) \text{ MeV}, \\ m_t &= [138_{-80}^{+52}(\text{expt.})_{-12}^{+20}(\text{Higgs})] \text{ GeV}, \\ \alpha_s(M_Z^2) &= 0.137 \pm 0.017(\text{expt.}) \pm 0.002(\text{Higgs}). \end{aligned} \quad (5.8)$$

The value of M_Z equals (within one GeV) the one obtained in the model independent fits.

Figure 5.3 shows the 68% confidence level contour (for the central value of M_H) in the $\alpha_s(M_Z^2)$ - m_t plane. The central values obtained for this Higgs mass and for the lower and upper limits are also given. The top mass obtained from the fit can be compared with the lower limit obtained in t -quark searches at the Tevatron by the CDF collaboration [65]: $m_t > 113$ GeV at 95% C.L. This is compatible with the value derived above.

The result for $\alpha_s(M_Z^2)$ is in agreement both with the value derived in the model independent approach and with the values obtained from the event shape analysis and the hadronic τ branching ratio measurement. Actually, the latter α_s determinations are more precise than the value (5.8). Combining them yields [29]

$$\alpha_s(M_Z^2) = 0.124 \pm 0.006. \quad (5.9)$$

Using this α_s value as an additional constraint in the fit one can now also obtain a more precise estimate of the top mass. The result found in a new fit ($\chi^2/N_{\text{DF}} = 87/103$) is

$$m_t = [160_{-54}^{+40}(\text{expt.})_{-15}^{+20}(\text{Higgs})] \text{ GeV}. \quad (5.10)$$

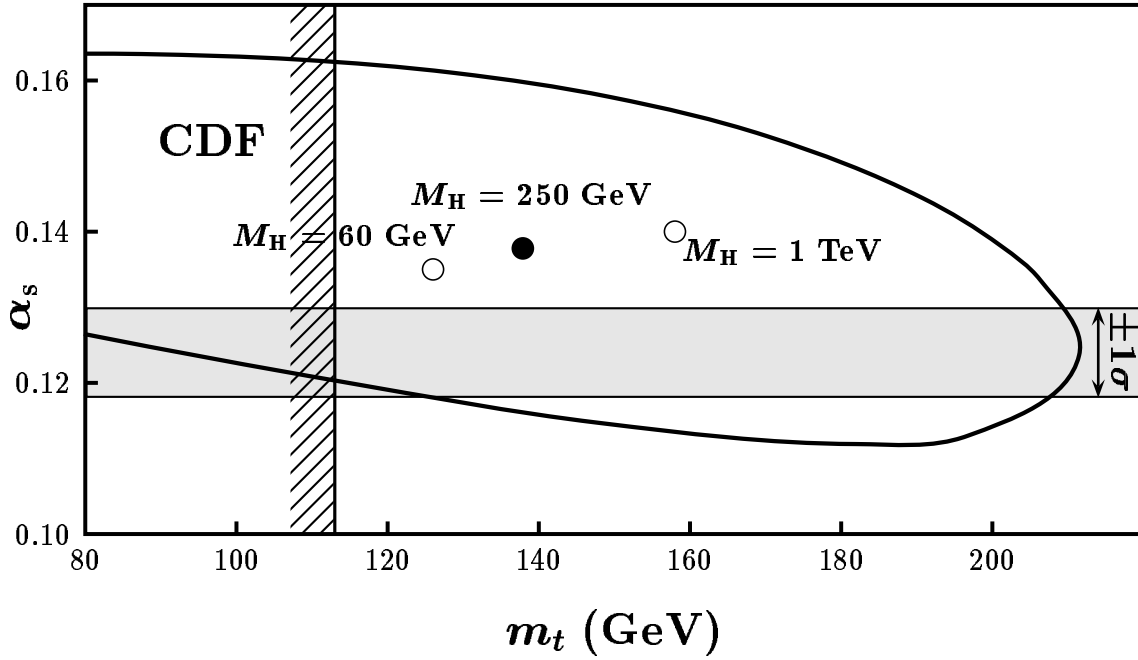


Figure 5.3: 68% Confidence Level contour in the $\alpha_s(M_Z^2)$ - m_t plane. Also shown is the lower limit obtained by CDF, $m_t > 113$ GeV at 95% C.L., and the $\pm 1\sigma$ -band for the $\alpha_s(M_Z^2)$ determinations following from the event shape analysis and the hadronic τ branching ratio measurement.

Since the four parameters mentioned above are the only free parameters of the Standard Model, all other parameters follow from them. In particular, for any given m_t and M_H , the value of Δr and hence $\sin^2 \theta_W$ can be calculated. Conversely, any additional information on $\sin^2 \theta_W$ can be used to improve the precision of the top mass determination.

For this purpose one can use the direct measurement $M_W/M_Z = 0.8798 \pm 0.0028$ by the CDF and UA2 collaborations [66]. This measurement implies a model independent determination of $\sin^2 \theta_W = 0.2260 \pm 0.0049$. Again, before imposing this extra information as an additional constraint on the fit, one has to verify that this value is consistent with the value derived from the unconstrained fit. Figure 5.4a shows that there is indeed a clear consistency, in the whole Higgs mass range considered. It is interesting to note that the precision obtained on $\sin^2 \theta_W$ from the lineshape data is comparable to that obtained by the CDF and UA2 collaborations.

Using the latter value as a constraint, again a good fit is obtained ($\chi^2/N_{\text{DF}} =$

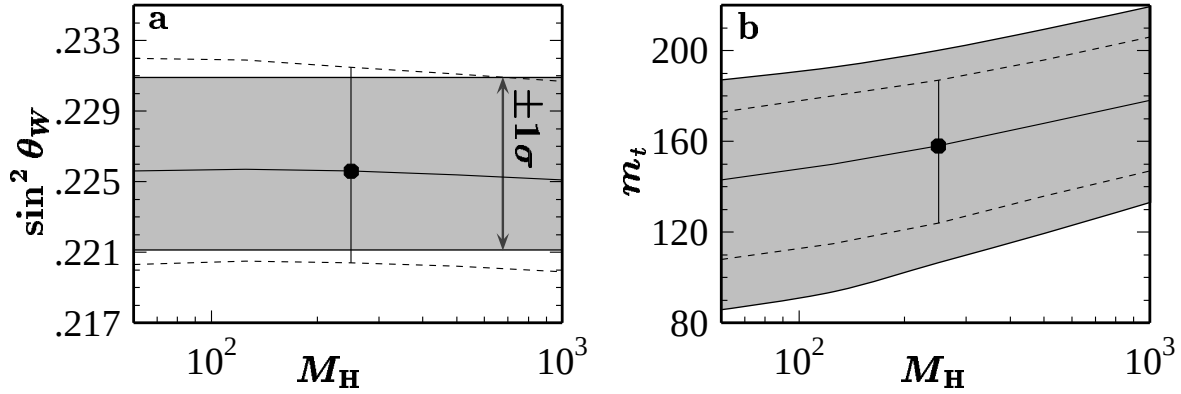


Figure 5.4: Variation of fitted quantities with Higgs mass. (a) variation of $\sin^2 \theta_W$ with M_H . Only the constraint on $\alpha_s(M_Z^2)$ is used; the constraint on $\sin^2 \theta_W$ is indicated by the gray band. (b) variation of m_t with M_H . The central line and the dashed lines denote the central value and the 68% C.L. band after imposing both the $\sin^2 \theta_W$ and the $\alpha_s(M_Z^2)$ constraints; the gray area represents the 68% C.L. band when only the constraint on $\alpha_s(M_Z^2)$ is imposed.

87/104) with the final result

$$m_t = [158^{+29}_{-34}(\text{expt.})^{+20}_{-15}(\text{Higgs})] \text{ GeV.} \quad (5.11)$$

The dependence of the fitted top mass on the input Higgs mass is displayed in figure 5.4b. As a reference, the values obtained without imposing the $\sin^2 \theta_W$ constraint are also shown.

Finally, from the top mass value (5.11) the following quantities can be derived:

$$\begin{aligned} \sin^2 \theta_W &= 0.2258 \pm 0.0037, \\ M_W &= (80.24 \pm 0.20) \text{ GeV.} \end{aligned} \quad (5.12)$$

As is to be expected from figure 5.4a, these values are practically independent of the Higgs mass assumed.

5.5 Conclusions

The model independent fits performed in section 5.1 demonstrate that using both the 1990 and 1991 cross sections and the leptonic forward-backward asymmetries it is possible to measure the mass and total decay width of the Z boson, M_Z and Γ_Z ,

with an accuracy of approximately 10 MeV. The small Standard Model dependence introduced in the fits in the form of assumptions for the interference terms affects exclusively M_Z and only by a few MeV. For the fitted Z mass, a good agreement with Standard Model predictions is obtained for all partial decay widths and effective coupling constants. An interpretation of the invisible width of the Z boson as originating from decays of the Z into light neutrinos with Standard Model couplings excludes the existence of four such neutrino species. Therefore at present there is no indication that the ‘minimal’ Standard Model does not describe Nature.

The processes observed have to be radiatively corrected. In the Standard Model context these corrections depend on the top mass. The sensitivity of the data to this dependence turns out to be large enough for the radiative corrections to yield a first estimate of the mass of the (as yet unobserved) top quark. It is not possible to obtain any information on the mass of the Higgs boson. The radiative corrections also allow an estimate of the $W^\pm$ mass with a precision comparable to the one obtained in direct measurements.

At present, the precision of the Z mass determination is limited by the fact that the LEP energy calibration through resonant depolarisation was achieved at one $\sqrt{s}$ point only. The energy points chosen for the 1993 data taking period are such that this calibration will be possible for all points. Combined with the increase in lineshape statistics for the 1993 running period, one therefore expects the determination of the Z mass to be substantially improved. Furthermore, the above improvements coupled with the reduction of the systematic uncertainties on the luminosity measurement are expected to significantly increase the precision of the total and partial decay widths. Lastly, the increased statistics of the on-peak 1992 and 1993 data for the leptonic forward-backward asymmetry measurements will also substantially improve the determination of the effective coupling constants. These facts in turn are expected to lead to a better precision of the $\alpha_s(M_Z^2)$ and m_t determinations. One can therefore expect that the data taken in the 1992 and 1993 running periods will allow an even more stringent test of the Standard Model.

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Summary

The subject of this thesis is the measurement of the total cross section for the process $e^+e^- \rightarrow \text{hadrons}$, on the Z resonance, as a function of the center of mass energy $\sqrt{s}$.

Within the context of the Standard Model, knowledge of the parameter M_Z entails a prediction with only a small theoretical uncertainty for other properties of the Z boson, e.g. its partial decay widths. Hence, one can test the model even though several other input parameters, notably the top mass m_t and the Higgs mass M_H , are still unknown. On the other hand, one can obtain information about these unknown parameters, by assuming the validity of the Standard Model and evaluating their effects on the so-called radiative corrections. At present, the accuracy of the theoretical predictions is such that – given sufficient *experimental* precision – both of these approaches give significant results.

A good experimental precision entails stringent requirements on the knowledge of both the center-of-mass energies and the measured cross sections. At the LEP collider at CERN, a precise absolute calibration (~ 6 MeV) of the former was achieved in 1991 for one energy point, using resonant depolarisation of the e^- beam. Other calibration techniques were used to extrapolate this result to the other energy points, albeit at the cost of some additional uncertainties.

The L₃ detector is used to study the interactions caused by colliding electron and positron beams. This detector is a multipurpose apparatus with emphasis on the detection of electrons, photons and muons. It contains a small tracking chamber, an electromagnetic calorimeter consisting of large arrays of BGO crystals, a hadronic calorimeter made out of a sandwich of uranium absorber plates and gas filled proportional wire chambers, and a large muon spectrometer comprising three layers of drift chambers. Scintillation counters ensure a precise event timing. The luminosity measurements are performed using two dedicated small-angle BGO calorimeters. The whole detector is housed in a solenoid providing a uniform magnetic field of 0.5 Tesla. A complex triggering and data acquisition system is responsible for an efficient data collection. Elaborate calibration and monitoring systems ensure a good detector precision.

The acquired data are passed through a set of complex event reconstruction programs: various subdetector objects are reconstructed, which are subsequently used to obtain global event information. The selection of hadronic events relies to a large extent on this global information. It is primarily based on the observed event

multiplicity and the (calibrated) total reconstructed event energy. Using Monte Carlo methods it is shown that on the basis of these features, it is possible to obtain a high selection efficiency ($\sim 99\%$) while maintaining a background contamination below 1%. Approximately 415,000 hadronic final states are selected from the 1990 and 1991 data.

The normalisations needed in order to translate the measured numbers of events into cross sections, the *luminosity measurements*, are performed using small angle Bhabha events. The main contributions to the systematic errors on the measured luminosities come from geometrical effects and from uncertainties in the theoretical cross sections; a normalisation uncertainty of 0.6% is obtained.

The high statistical accuracy obtained allows the study of various systematic effects. It is shown that the trigger system can be regarded as fully efficient. The measured cross sections are corrected for the natural spread of the LEP beam energy. The fill-to-fill cross sections show only the fluctuations to be expected on the basis of statistical considerations and no significant variations with time. The subtraction of a ‘non-resonant’ background due to two-photon processes and beam-gas events leads to an uncertainty of 30 (20) pb for the 1990 (1991) data. When varying the selection criteria, the measured cross sections are stable to within 0.10% (0.05%). The effect of a random background is estimated and yields $\sqrt{s}$ dependent uncertainties of only a few per mille. The Monte Carlo modelling of the fragmentation process is shown to contribute a 0.15% uncertainty.

The measured hadronic cross sections, together with leptonic cross section and forward-backward asymmetry data, are fitted to theoretical predictions. The mass, total decay width and hadronic decay width of the Z boson are measured to be $(91197 \pm 6 \pm 7)$ MeV, $(2492 \pm 10 \pm 5)$ MeV and (1753 ± 14) MeV respectively, where the first error represents the experimental uncertainty, and the second error (where present) is due to the uncertainty in the LEP beam energy values. Good fits are obtained. The measured partial decay widths as well as the effective coupling constants are in good agreement with Standard Model predictions. The number of light neutrino generations inferred from the total and partial decay widths, $N_\nu = 2.98 \pm 0.06$, clearly excludes any integer number other than three. The value of the strong coupling constant α_s deduced from the ratio of the hadronic to leptonic partial decay widths, $\alpha_s(M_Z^2) = 0.148_{-0.024}^{+0.012}$, is consistent with other, independent measurements of this quantity.

Interpreted in the Standard Model context, the data are used to obtain information on the top mass. Constraining both $\alpha_s(M_Z^2)$ and $\sin^2 \theta_W = 1 - M_W^2/M_Z^2$ to other independent measurements, one obtains the most accurate results: $m_t = [158_{-34}^{+29}(\text{expt.})_{-15}^{+20}(\text{Higgs})]$ GeV, where the second error represents the variation of the fitted m_t when varying M_H between 60 and 1000 GeV. The top mass found leads to a precise, though indirect, determination of the $W^\pm$ mass: $M_W = (80.24 \pm 0.20)$ GeV.

Samenvatting

Meting van Hadronische Werkzame Doorsnedes op de Z-resonantie met de L₃-detector

Het onderwerp van dit proefschrift is de meting van de totale werkzame doorsnede voor het proces $e^+e^- \rightarrow \text{hadronen}$, op de Z-resonantie, als functie van de zwaartepunts-energie $\sqrt{s}$.

Binnen de context van het Standaard Model voor electrozwakke interacties is de massa M_Z van het Z-boson tot op grote hoogte bepalend voor vele andere eigenschappen van dit deeltje. Dit impliceert dat metingen van M_Z gebruikt kunnen worden voor tests van het Standaard Model zelfs als niet alle parameters van het model bekend zijn. Aan de andere kant kan men, aannemend dat het model inderdaad geldig is, deze onbekende parameters (in het bijzonder de massa van de top quark m_t en de Higgs massa M_H) proberen te bepalen; daarbij wordt gebruik gemaakt van de grootte van de zogenaamde stralingscorrecties. De theoretische voorspellingen zijn voldoende precies om, indien de benodigde metingen maar accuraat genoeg zijn, beide invalshoeken te kunnen volgen.

Een goede experimentele precisie vergt nauwkeurige metingen zowel van de zwaartepunts-energieën als van de corresponderende werkzame doorsnedes. Bij de LEP-botsingsring op CERN werd in 1991 één van voornoemde energiepunten, door middel van resonante depolarisatie van de e^- bundel, tot op ongeveer 6 MeV gecalibreerd. Met behulp van andere calibratietechnieken werden vervolgens, zij het ten koste van enige extra onzekerheden, ook de andere energiepunten gecalibreerd.

De interacties, veroorzaakt door de botsende elektronen- en positronenbundels, zijn bestudeerd met behulp van de L₃-detector. Dit is een multifunctionele opstelling die geoptimaliseerd is voor de detectie van elektronen, fotonen en muonen. Ze omvat een kleine sporenkamer, een electromagnetische calorimeter die bestaat uit rijen BGO-kristallen, een hadronische calorimeter die opgebouwd is uit lagen van uranium platen met daartussen gasgevulde dradenkamers, en een grote muon spectrometer die bestaat uit drie lagen driftkamers. Scintillatietellers zorgen voor

een preciese timing van de interacties. Twee kleine BGO-calorimeters, geplaatst onder kleine hoeken met de LEP bundelpijp, verrichten de meting van de luminositeit. De hele detector-opstelling is geplaatst in een uniform, axiaal magnetisch veld van 0.5 Tesla gegenereerd door een solenoïde magneet. Een complex beslissings- en data-acquisitie systeem zorgt ervoor dat de data op een efficiënte manier geregistreerd worden. De vereiste detector-precisie wordt gewaarborgd door uitgebreide bewakings- en calibratiesystemen.

De geregistreeerde data worden gereconstrueerd met behulp van een reeks complexe software-programma's. Deze maken in eerste instantie subdetector-'objecten', die vervolgens gecombineerd worden tot globale informatie. De selectie van de hadronische data is voor een groot deel geënt op deze globale informatie; ze is voornamelijk gebaseerd op de gemeten multipliciteiten en totale energieën. Door middel van Monte Carlo methoden kan aangetoond worden dat het met behulp van deze kenmerken mogelijk is hadronen te selecteren met hoge efficiëntie (ongeveer 99%), terwijl de geselecteerde achtergrond beneden 1% blijft. Op deze manier zijn uit de 1990 en 1991 data ongeveer 415.000 hadronische eindtoestanden geselecteerd.

De meting van de luminositeiten, dat wil zeggen van de normalisaties die nodig zijn om gemeten aantallen hadronische eindtoestanden te vertalen in werkzame doorsnedes, geschiedt met behulp van de sterk voorwaarts gerichte Bhabha eindtoestanden. De voornaamste systematische onzekerheden in de gemeten luminositeiten worden veroorzaakt door geometrische effecten en onzekerheden in de theoretische beschrijving van het Bhabha-proces; de uiteindelijke onzekerheid op de normalisatie bedraagt 0.6%.

Door de grote hoeveelheid hadronische data is het mogelijk een grote verscheidenheid aan systematische effecten te onderzoeken. Zo kan aangetoond worden dat, voor het doeleinde van dit proefschrift, het beslissingssysteem om interacties al dan niet te registreren als volledig efficiënt beschouwd kan worden. De werkzame doorsnedes worden gecorrigeerd voor de natuurlijke spreiding van de bundelenergieën; wanneer deze doorsnedes opgesplitst worden in kleinere tijdseenheden, blijken hun variaties compatibel met statistische fluctuaties en worden er geen significante tijdsafhankelijkheden gesobserveerd. Het aftrekken van een 'niet-resonante' achtergrond brengt een onzekerheid van 30 pb (20 pb) met zich mee voor de 1990 (1991) data. Wanneer de selectiecriteria gevarieerd worden veranderen de gemeten werkzame doorsnedes minder dan 0.10% (0.05%) voor de 1990 (1991) data. De verwachting is dat de data ook een 'toevals'-achtergrond bevatten; correctie voor deze achtergrond leidt tot een onzekerheid van enkele promilles. Het verschil in keuze van het model voor het fragmentatieproces zoals geïmplementeerd in Monte Carlo programma's resulteert in een 0.15% onzekerheid op de gemeten hadronische werkzame doorsnedes.

Aan deze gemeten hadronische werkzame doorsnedes, samen met de resultaten voor de leptonische totale werkzame doorsnedes en voorwaarts-achterwaarts asymmetrieën, worden theoretische voorspellingen aangepast. De massa, de totale vervalsbreedte en de hadronische vervalsbreedte van het Z-boson worden respectievelijk

bepaald als $M_Z = (91197 \pm 6 \pm 7) \text{ MeV}$, $\Gamma_Z = (2492 \pm 10 \pm 5) \text{ MeV}$ en $\Gamma_{\text{had}} = (1753 \pm 14) \text{ MeV}$; in deze resultaten correspondeert de eerste fout met de experimentele onzekerheid, en de tweede (indien aanwezig) met de onzekerheden in de LEP energiecalibratie. De model-onafhankelijke parametrisaties sluiten goed aan bij de data. Zowel de partiële en totale vervalsbreedtes als de effectieve koppelingsconstanten zijn in goede overeenkomst met voorspellingen van het Standaard Model. Het aantal lichte neutrino-generaties zoals bepaald uit de totale en partiële vervalsbreedtes van het Z-boson, $N_\nu = 2.98 \pm 0.06$, sluit ieder geheel getal anders dan drie uit. De waarde van de sterke koppelingsconstante $\alpha_s(M_Z^2)$, voortvloeiend uit de verhouding van de hadronische tot de leptonische vervalsbreedte, $\alpha_s(M_Z^2) = 0.148^{+0.012}_{-0.024}$, is consistent met andere, onafhankelijke metingen van deze parameter.

De data worden, binnen de context van het Standaard Model, gebruikt om een schatting van de top-massa te verkrijgen. Door combinatie van deze data met bovengenoemde onafhankelijke metingen van $\alpha_s(M_Z^2)$ en metingen van $\sin^2 \theta_W = 1 - M_W^2/M_Z^2$, wordt de nauwkeurigste waarde gevonden: $m_t = [158^{+29}_{-34}(\text{expt.})^{+20}_{-15}(\text{Higgs})] \text{ GeV}$. De eerste fout beschrijft de onzekerheid in m_t voor een centraal genomen waarde $M_H = 250 \text{ GeV}$ voor de (nog) onbekende massa van het Higgs-deeltje, en de tweede geeft de variatie aan in de aldus bepaalde top-massa wanneer M_H gevarieerd wordt tussen 60 en 1000 GeV. Het gevonden resultaat voor m_t leidt tot een preciese, zij het indirecte, bepaling van de $W^\pm$ massa: $M_W = (80.24 \pm 0.20) \text{ GeV}$.

Curriculum Vitae

De schrijver van dit proefschrift werd geboren op 16 juli 1967 te 's-Hertogenbosch. Vanaf 1979 bezocht hij het Maurick College te Vught, wat in 1985 resulteerde in het behalen van het VWO-diploma.

In 1985 begon hij met de studie Natuurkunde aan de Katholieke Universiteit Nijmegen. In 1986 behaalde hij hier zijn propaedeuse. In 1988 begon hij met afdelingswerk op de afdeling Theoretische Hoge Energie Fysica. Onder begeleiding van prof. dr. C. Dullemond werkte hij aan zijn afstudeerscriptie, getiteld "Renormalization in Anti-de Sitter space-time". In november 1989 werd hem door de faculteit der Natuurwetenschappen het doctoraal diploma uitgereikt.

In januari 1990 begon hij zijn aanstelling als Onderzoeker In Opleiding (O.I.O.) in dienst van de stichting Fundamenteel Onderzoek der Materie (FOM), verbonden aan de vakgroep Experimentele Hoge Energie Fysica van bovengenoemde universiteit. In het kader van zijn aanstelling was hij van juni 1990 tot mei 1992 werkzaam bij het L_3 experiment, uitgevoerd bij het CERN te Genève. Hij leverde daar een contributie aan dit experiment en begon met het onderzoek dat leidde tot dit proefschrift.

De laatste anderhalf jaar van zijn aanstelling was hij wederom werkzaam in Nijmegen. Hier schreef hij zijn proefschrift onder begeleiding van prof. dr. R.T. Van de Walle en prof. dr. E.W. Kittel. In juli 1992 gaf hij namens het L_3 -samenwerkingsverband op het "XXII International Symposium on Multiparticle Dynamics" te Santiago de Compostela (Spanje) een voordracht over "Inclusive Hadron Production in L_3 ". Ook was hij betrokken bij het geven van onderwijs aan 2^e- en 3^e-jaars natuurkunde studenten.