SEARCHES FOR PAIR-PRODUCED CHARGED HIGGS BOSONS AND
NEUTRAL SUPERSYMMETRIC HIGGS BOSONS IN Z° DECAYS

Douglas F. Cowen

Under the supervision of Professor Sau Lan Wu
at the University of Wisconsin-Madison

ABSTRACT

Searches for pair-produced charged Higgs bosons and neutral supersymmetric
Higgs boson in decays of the Z° have been performed using the ALEPH detector
at LEP for the decay channels $H^+ H^- \rightarrow c\bar{s}$, $H^+ H^- \rightarrow c\bar{b}$, and $h^0 A^0 \rightarrow b\bar{b} b\bar{b}$. With 1.17 pb$^{-1}$ of integrated luminosity, corresponding to about 25000
hadronic decays of the $Z^0$, the charged Higgs has been excluded at 95% CL in
the mass range 14.5–35.2 GeV for $\text{BR}[H^\pm \rightarrow cs] = 100\%$ and 16.4–35.4 GeV for
$\text{BR}[H^\pm \rightarrow cb] = 100\%$. The neutral supersymmetric Higgs has been excluded
at 95% CL in the mass range 15.0–40.0 GeV for $M_{h^0} = M_{A^0}$, and as high as
$M_{A^0} = 44.0$ GeV for $M_{h^0} = 24.0$ GeV.
Acknowledgements

There are two broad categories of people whom I would like to thank: those who helped me, in one way or another, to do the work this thesis represents, and those who patiently waited for me to finish it. Clearly, these two categories are substantially overlapping. I will start with the former.

It is not possible in a few sentences to thank my advisor, Sau Lan Wu, for all that she has done for me in the six years I have worked with her as a graduate student. Her unswerving dedication to physics and relentless drive and enthusiasm are exemplary. I am indebted to the past and present Wisconsin group for answering questions and generally making the working environment at CERN a pleasant one. In particular, I would like to thank Mike Mermikides and Steve Ritz, who got me started on the charged Higgs analysis, and Leo Bellantoni and Jo Pater, who both did a lot of work with the charged Higgs in addition to my own efforts. My officemate Robert Johnson deserves my gratitude for his forbearance, and for stoically accepting the fact that he was often the first person I turned to when I had a question. Olga Mermikides and Tricia Van Gene, the former and present secretaries of the Wisconsin group, have both
been very helpful in all matters administrative. Last but not least, those people who proofread this thesis have my sincere gratitude.

Outside of the Wisconsin group, I would like to thank T.T. Wu and Herbi Dreiner for their careful reading of the first chapter. The neutral SUSY Higgs analysis was done with the help of two French collaborators, Jean-Francois Grivaz and Patrick Janot. I thank Julie Kreunen for her initial support and for the literary reference. I am indebted to Dick Jared for what he taught me about hardware, and to Ron Settles for helping me while I was at MPI in Munich. Finally, I offer my thanks to Ugo Camerini for his help in organizing the defense of this thesis.

It has been my great fortune to have made many good friends during my stay at CERN. Of these, I would especially like to thank Jenny Thomas, for her moral support, and Dave Muller, Steve Ritz, and Jim Wear who as housemates often helped smooth over the rough times.

Madame Bramerel me manquera beaucoup, et j’espère que je pourrai lui rendre visite encore plusieurs fois. Elle m’a dit la vérité en disant “Tu auras des racines ici, Doug.”

None of my accomplishments would have been possible without the endless support and encouragement of my family, whose patience with me applies not only to the time it took to complete this thesis, but to all other aspects of my upbringing. My wife Clare is the best proof that the two categories above are overlapping. Her patience, peppered with occasional bursts of impatience, and her emotional and moral support for nearly three years have served to keep
me going in this difficult endeavor. It would have been extraordinarily trying without her.
Contents

Abstract i

Acknowledgments ii

1 Introduction 1
   1.1 Overview ........................................ 2
   1.2 Fundamental Particles and Forces .............. 3
   1.3 Yang–Mills Theory .............................. 7
   1.4 Spontaneously Broken Symmetries ............... 8
   1.5 The Standard Model ............................. 12
   1.6 Non-Minimal Higgs Sectors .................... 14
   1.7 Properties of the Higgs Boson ................. 15
      1.7.1 Higgs Production Mechanisms ............ 15
      1.7.2 Higgs Decay Mechanisms ................ 18
      1.7.3 Current Experimental Limits ............. 19
2 The LEP Storage Ring and the ALEPH Detector .............................. 20
  2.1 The LEP Storage Ring ............................................. 20
  2.2 The ALEPH Detector .................................................. 22
    2.2.1 Introduction ................................................. 24
    2.2.2 The ITC ..................................................... 26
    2.2.3 The TPC ..................................................... 26
    2.2.4 The ECAL .................................................. 32
    2.2.5 The Superconducting Solenoid ............................... 35
    2.2.6 The HCAL and the Muon Chambers ..................... 36
    2.2.7 The Luminosity Monitors .................................. 40

3 Event Triggering, Reconstruction, and Simulation .................. 41
  3.1 Event Triggering and Reconstruction ............................. 42
  3.2 Event Simulation ................................................ 44
    3.2.1 $Z^0 \rightarrow q\bar{q}(g)$ Background Simulation ........ 45
    3.2.2 Charged and Neutral Higgs Signal Simulation .......... 48

4 The $H^+H^-$ and $A^0 h^0$ Four-Jet Analyses .................. 50
  4.1 Charged Higgs Selection Criteria .................................. 61
  4.2 Neutral SUSY Higgs Selection Criteria .......................... 74
    4.2.1 Electron Identification ..................................... 77
    4.2.2 Muon Identification ....................................... 81
    4.2.3 Selection Criteria Based on Lepton Identification .... 85
5 Experimental Results
  5.1 Systematics ............................................. 93
  5.2 Calculating the Limits ................................. 95
    5.2.1 Limits on Charged Higgs Production ............ 96
    5.2.2 Limits on Neutral SUSY Higgs Production ....... 107
  5.3 Conclusion ............................................. 111

A The TPC Gating System ..................................... 113
  A.1 Introduction ........................................... 113
  A.2 Principle of Gating .................................... 113
  A.3 Design Considerations ................................. 115
  A.4 Performance ........................................... 116

B Compensation for Edge Effects in TPC Sectors .......... 123
  B.1 Introduction ........................................... 123
  B.2 Apparatus ............................................. 125
  B.3 Procedure ............................................. 127
  B.4 Results ................................................. 129
  B.5 Conclusion ............................................. 132
Chapter 1

Introduction

Significant increases in the energy delivered by particle accelerators have often coincided with the discovery of new particles. Many of these discoveries have occurred at experiments at colliding beam facilities using either $p\bar{p}$ or $e^+e^-$ beams. With the advent of LEP at CERN, an unprecedented number of $Z^0$ bosons are being produced. Although the $Z^0$ boson was discovered at a $p\bar{p}$ machine [1], the large, high-purity sample being collected by the LEP experiments permits a much closer scrutiny of the $Z^0$'s properties and the properties of its decay products. In either of these arenas new, unexpected properties of matter can potentially be discovered.

This thesis presents a search for the $Z^0$ decay in which scalar particles, known as Higgs bosons [2], are pair-produced. The searches are made for the case where these Higgs particles each decay into a pair of quarks, which then individually fragment into many particles. These particles are collimated into "jets" by virtue of the initial momentum of their parent partons. Thus, the signature left by these Higgs bosons takes the form of a multi-jet, high-multiplicity event. Events with
such a topology are, in fact, numerous at LEP due to the high cross section for $Z^0 \rightarrow q\bar{q}(g)$ processes, which also have a multi-jet signature. Hence, many different selection criteria, each taking advantage of an aspect of the Higgs boson signature which is different from that of the $Z^0 \rightarrow q\bar{q}(g)$ signature, are applied to the data in an effort to enhance the potential Higgs signal to the point where it would be measurably distinct from the $Z^0 \rightarrow q\bar{q}(g)$ background.

The remainder of this chapter is concerned with the physical theory which anticipates the existence of Higgs bosons, and with its explicit predictions regarding their production and decay. Past experimental limits are also discussed. Chapter 2 presents the LEP storage ring and the ALEPH detector, with a description of the subcomponents and performance of the latter. In Chapter 3 the triggering of the detector, the reconstruction of the data and simulated events, and the Monte Carlo simulations are described. Chapter 4 provides a detailed picture of the analyses used to set the limits on pair-produced Higgs bosons, the procedure for which is presented in Chapter 5. This last chapter also contains a discussion of the systematic errors and the conclusion.

1.1 Overview

Bringing together many years of work by many physicists, Glashow, Weinberg and Salam proposed in 1967 a model in which the weak and electromagnetic forces were unified into one force, the "electroweak" force. This model has been so successful—no measurement to date runs counter to its predictions—that, in conjunction with Quantum Chromodynamics (QCD), it has come to be called
the "Standard Model" [3]. Electroweak theory's most prominent achievements have been the prediction of the properties of the intermediate vector bosons and its incorporation of weak neutral currents. Yet, in spite of its great successes, the theory is considered incomplete. Chief amongst its shortcomings are that it makes no predictions regarding the number of fermion generations, the quantization of charge, or the values of the fermion masses. Also, it predicts the existence of the "Higgs" boson [2], or, in extensions of the model, several such bosons. The Higgs mechanism, by which these bosons are given mass, is a central element of electroweak theory but such bosons have not yet been observed, nor does the model make any predictions regarding their masses. In this thesis, the case where the Standard Model is minimally extended to have five Higgs bosons is studied.

Electroweak theory arises conceptually from the interplay of Yang–Mills theories and spontaneous symmetry breaking. The theory of the electroweak force will be discussed in this context later on in this chapter. To provide the proper backdrop the particle phenomenology and the forces of the model are discussed first.

1.2 Fundamental Particles and Forces

The fundamental, point-like particles of the Standard Model consist of leptons and quarks. The former do not participate in the strong interaction while the latter do. Quarks form heavier composite states known as hadrons. The model
<table>
<thead>
<tr>
<th>Particle Name</th>
<th>Symbol</th>
<th>Charge</th>
<th>Mass (MeV)</th>
<th>Lifetime(s)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Electron Neutrino</td>
<td>$\nu_e, \bar{\nu}_e$</td>
<td>0</td>
<td>&lt; 0.018</td>
<td>$\sim \infty$</td>
</tr>
<tr>
<td>Muon Neutrino</td>
<td>$\nu_\mu, \bar{\nu}_\mu$</td>
<td>0</td>
<td>&lt; 0.25</td>
<td>$\sim \infty$</td>
</tr>
<tr>
<td>Tau Neutrino*</td>
<td>$\nu_\tau, \bar{\nu}_\tau$</td>
<td>0</td>
<td>&lt; 35.</td>
<td>$\sim \infty$</td>
</tr>
<tr>
<td>Electron</td>
<td>$e^\pm$</td>
<td>$\pm e$</td>
<td>0.511</td>
<td>$\sim \infty$</td>
</tr>
<tr>
<td>Muon</td>
<td>$\mu^\pm$</td>
<td>$\pm e$</td>
<td>105.66</td>
<td>$2.19 \times 10^{-6}$</td>
</tr>
<tr>
<td>Tau</td>
<td>$\tau^\pm$</td>
<td>$\pm e$</td>
<td>1784</td>
<td>$3.04 \times 10^{-13}$</td>
</tr>
</tbody>
</table>

Table 1.1: The leptons of the Standard Model. *The $\nu_\tau$ has not yet been observed, but there is strong circumstantial evidence that it exists (see Ref. 4).

also provides a framework for understanding three of the four known fundamental interactions, or forces: weak, electromagnetic and strong. The theory of the strong force, while similar in structure to that of the electroweak, is probably an exact symmetry and as such does not require spontaneous symmetry breaking. The fourth force, gravity, has not been successfully described as a quantum field theory. For each of these forces there is an associated boson(s) which transmits the force between the fundamental particles.

The leptons are listed in Table 1.1. The hadrons are numerous but their spectrum is well understood in terms of their constituent building blocks, the quarks. The hadrons can be split into two broad categories, mesons and baryons, consisting of a bound state of two and three quarks, respectively. Although free quarks have never been observed in the laboratory a large enough body of circumstantial evidence has been accumulated over the years so that physicists
Table 1.2: The quarks of the Standard Model. *Since quarks do not exist as free particles, their masses cannot be directly measured. †The top quark has not yet been observed.

```
<table>
<thead>
<tr>
<th>Quark</th>
<th>Symbol</th>
<th>Charge</th>
<th>Mass* (GeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Up</td>
<td>u</td>
<td>$+\frac{2}{3}e$</td>
<td>0.006</td>
</tr>
<tr>
<td>Down</td>
<td>d</td>
<td>$-\frac{1}{3}e$</td>
<td>0.01</td>
</tr>
<tr>
<td>Strange</td>
<td>s</td>
<td>$-\frac{1}{3}e$</td>
<td>0.19</td>
</tr>
<tr>
<td>Charm</td>
<td>c</td>
<td>$+\frac{2}{3}e$</td>
<td>1.35</td>
</tr>
<tr>
<td>Bottom</td>
<td>b</td>
<td>$-\frac{1}{3}e$</td>
<td>$\approx 5$</td>
</tr>
<tr>
<td>Top†</td>
<td>t</td>
<td>$+\frac{2}{3}e$</td>
<td>?</td>
</tr>
</tbody>
</table>
```

The fundamental particles which transmit the electroweak force are the photon and the three intermediate vector bosons, the $W^{\pm}$ and the $Z^0$. The strong force has eight exchange particles called gluons, and gravity is presumed to be transmitted by a spin-2 particle called the graviton. The photon was first observed via the photoelectric effect some time ago [5]. Gluons were first observed in 1979 [6], and the intermediate vector bosons were first observed in 1983 [1]. There is no experimental evidence for the graviton or gravitational waves. The properties of these exchange particles are listed in Table 1.3.
<table>
<thead>
<tr>
<th></th>
<th>Gravitational</th>
<th>Weak</th>
<th>Electromagnetic</th>
<th>Strong</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Range</strong></td>
<td>$\infty$</td>
<td>$\ll 10^{-14}$ cm</td>
<td>$\infty$</td>
<td>$10^{-13} - 10^{-14}$ cm</td>
</tr>
<tr>
<td><strong>Strength</strong></td>
<td>$5.9 \times 10^{-39}$</td>
<td>$1.02 \times 10^{-5}$</td>
<td>$\alpha_{em} = \frac{1}{137}$</td>
<td>$\alpha_{strong} \approx 0.1$</td>
</tr>
<tr>
<td><strong>Particles acted upon</strong></td>
<td>All</td>
<td>Hadrons and Leptons</td>
<td>Charged particles</td>
<td>Hadrons</td>
</tr>
<tr>
<td><strong>Particles exchanged</strong></td>
<td>Gravitons(?)</td>
<td>Intermediate vector bosons</td>
<td>Photons</td>
<td>Gluons</td>
</tr>
<tr>
<td><strong>Mass of exchanged particle</strong></td>
<td>0</td>
<td>$80 - 100$ GeV</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td><strong>Spin of exchanged particle</strong></td>
<td>2</td>
<td>1</td>
<td>1</td>
<td>1</td>
</tr>
</tbody>
</table>

Table 1.3: The four fundamental forces.
1.3 Yang–Mills Theory

Yang–Mills theory, one of the mainstays of electroweak theory, is a generalization of the quantum mechanical gauge invariance principle. This principle states that quantum mechanical observables are unaffected by gauge or phase transformations of the form

$$\Psi(x) \rightarrow \Psi'(x) = e^{i\theta} \Psi(x).$$

(1.1)

This gauge transformation is known as a global phase transformation because $\theta$ is independent of $x$. It is also an Abelian transformation because successive transformations commute:

$$e^{i\theta_1} e^{i\theta_2} = e^{i\theta_2} e^{i\theta_1}.$$  

(1.2)

In the first step toward constructing a Yang–Mills theory, one demands that the phase invariance be local instead of global, i.e., that $\theta \rightarrow \theta(x)$ in Eq. 1.1 so that one can make a phase transformation independently at every point in spacetime. Yang and Mills eloquently stated the reason for insisting on local gauge invariance in their seminal 1954 paper [21]:

[In the absence of the electromagnetic interaction,] once one chooses what to call a proton, what a neutron, at one space-time point, one is then not free to make any choices at other space-time points...It seems that this is not consistent with the localized field concept that underlies the usual physical theories.
When Yang and Mills did the work culminating in their 1954 paper, they were motivated by a desire to create a theory of nucleons which conserved hadronic isospin. For this reason they went a step further and studied the consequences of a gauge transformation which was both local and non-Abelian with an underlying $SU(2)$ structure. They constructed a theory which had self-interacting vector bosons as the mediators of the force between nucleons. In spite of the mathematical beauty of the theory, the internucleon force operates only at short range. Hence, the bosons responsible for it would have to be massive, but the theory could not predict anything about the mass of these particles. One requires the bosons of the electroweak force to be massive for similar reasons.

1.4 Spontaneously Broken Symmetries

It was not until one decade after the introduction of Yang–Mills theory that Higgs and Englert and Brout [2] discovered the other mainstay of modern electroweak theory in showing how one could give mass to the massless vector bosons of the Yang–Mills theory. In essence, the idea is to keep the Lagrangian invariant under certain symmetry transformations while relaxing this requirement for the lowest energy state, i.e., the physical vacuum. Such symmetries are known as “spontaneously broken symmetries” and they give rise to massless scalar bosons, called Goldstone bosons [8]. However, an interaction between the massless Goldstone bosons and the massless vector bosons gives the latter mass and simultaneously entirely removes the former from the particle spectrum. This interplay is known as the Higgs mechanism.
Consider the Lagrangian for a complex scalar field \( \phi \), representing a particle and its antiparticle, each having the same mass \( \mu \):

\[
\mathcal{L} = (\partial_{\mu} \phi^\dagger)(\partial^\mu \phi) - V(\phi),
\]

\[
V(\phi) = \mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2. \tag{1.3}
\]

This particular potential, the first two terms in a general expansion of \( V \) in powers of \( \phi \), is chosen because it is invariant under the global transformation of Eq. 1.1. Furthermore, this choice of potential provides mass for the particles via its first term, and allows the fields to self-interact through its second term. Finally, higher order terms in the expansion are non-renormalizable and would render the theory unphysical in the framework of perturbation theory.

Depending on what one chooses for the sign of \( \mu^2 \) in Eq. 1.3, one gets either a manifest symmetry in which the physical vacuum obeys the original symmetry, or spontaneous symmetry breaking in which it does not. A plot of \( V(\phi) \) vs. \( \phi \) for \( \mu^2 > 0 \) reveals that the potential has a classical minimum at

\[
\phi \equiv \langle \phi \rangle_0 = 0.
\]

However, making the same plot for \( \mu^2 < 0 \) reveals that there are two such minima,

\[
\phi \equiv \langle \phi \rangle_0 = \pm \sqrt{-\frac{\mu^2}{2\lambda}}, \tag{1.4}
\]

where the arbitrary phase has been set to zero. These two behaviors are shown in Fig. 1.1.

Writing \( \phi \) as \( \phi = \phi_R + i\phi_I \), the Lagrangian is
Figure 1.1: The Potential Energy for the Scalar Field $\phi$. 

$V(|\phi|)$

$\mu^2 > 0$

$\mu^2 < 0$
\[ \mathcal{L} = \frac{1}{2} \left( \partial_\mu \phi_R \partial^\mu \phi_R + \partial_\mu \phi_I \partial^\mu \phi_I \right) - \mu^2 \left( \frac{\phi_R^2 + \phi_I^2}{2} \right) - \lambda \left( \frac{\phi_R^2 + \phi_I^2}{2} \right)^2 \] (1.5)

For \( \mu^2 > 0 \), the two scalar fields \( \phi_R, \phi_I \) each have mass \( \mu \), and they interact via the third term of the Lagrangian 1.5. In contrast, a very different situation obtains for \( \mu^2 < 0 \). In this case perturbation theory leads one to describe the fields as fluctuations about one of the minima of \( V(\phi) \):

\[ \phi = \frac{v + \rho(x)}{\sqrt{2}} \exp i \chi(x)/v. \] (1.6)

where

\[ \frac{v}{\sqrt{2}} = \langle \phi \rangle_0 \left( = \sqrt{-\frac{\mu^2}{2\lambda}} \right). \] (1.7)

At this point one has singled out one of the minima, \( i.e., \) one has chosen a particular phase for \( \langle \phi \rangle \), and so the vacuum no longer obeys the original gauge symmetry—the symmetry has been spontaneously broken. Writing the derivative terms of the Lagrangian 1.5 in terms of Eq. 1.6, one finds that the \( \rho \)-field has a mass \( \mu \) while the \( \chi \)-field, with no similar corresponding terms in the potential, remains massless. Nambu and Goldstone [8,9] showed that such a massless scalar particle will always be produced when one spontaneously breaks a continuous symmetry. These particles have since been known as Nambu-Goldstone bosons. However, the Nambu-Goldstone theorem can be violated, and the scalar particles can acquire mass, in the presence of gauge fields like the electromagnetic field \( A^\mu \). Incorporating the electromagnetic field \( A^\mu \) in the Lagrangian, and then
spontaneously breaking the symmetry, results in a massive scalar field and a massive vector field $A'^\mu$. The new massive scalar field is the Higgs field and its associated boson is called the Higgs boson. This procedure—in which the massless Nambu-Goldstone bosons are removed from the spectrum, the gauge field is given mass, and the Higgs boson is produced—is known as the Higgs mechanism [2].

1.5 The Standard Model

The underlying symmetry of electroweak theory is $SU(2)_L \otimes U(1)_Y$. This symmetry has four generators $T_1, T_2, T_3,$ and $Y$ and hence four gauge bosons. The two bosons associated with $T_3, Y$ are neutral, and since the linear combination

$$T_3 + \frac{Y}{2}$$  \hspace{1cm} (1.8)

leaves the vacuum invariant after spontaneous symmetry breaking, it is associated with the massless photon. The three remaining bosons acquire mass via the Higgs mechanism.

The covariant derivative appropriate for electroweak $SU(2)_L \otimes U(1)_Y$ symmetry is

$$D_\mu = \partial_\mu + igW \cdot T + ig'\frac{Y}{2}B_\mu$$  \hspace{1cm} (1.9)

\footnote{The "L" subscript on $SU(2)_L$ is used to indicate that experimentally only the left handed components of fermions have been observed to partake in the weak interactions, and the "Y" subscript on $U(1)_Y$ is used to indicate that the associated generator is weak hypercharge and not electromagnetic charge.}
where the $g$ term represents the $SU(2)$ interaction, the $g'$ term the $U(1)$ interaction, and $T = \tau/2$, where $\tau$ are the familiar Pauli spin matrices. Introducing the $SU(2)$ doublet

$$
\phi = \begin{bmatrix} \phi^+ \\
\phi^0 \end{bmatrix}
$$

(1.10)

into the Lagrangian and spontaneously breaking the symmetry, one obtains the vacuum expectation value

$$
\langle \phi \rangle_0 = \begin{bmatrix} 0 \\
v/\sqrt{2} \end{bmatrix}.
$$

(1.11)

In the Lagrangian the following quantity provides the masses for the three remaining vector gauge bosons (the $W^\pm$ and the $Z^0$):

$$
D_\mu \langle \phi \rangle^\dagger D_\mu \langle \phi \rangle = \left| (igW_\mu \cdot T + ig'\frac{Y}{2}B_\mu) \begin{bmatrix} 0 \\
v/\sqrt{2} \end{bmatrix} \right|^2 = \frac{iv}{2\sqrt{2}} \left| g(W_{1\mu} - iW_{2\mu}) \right|^2
$$

(1.12)

where $Y\langle \phi \rangle_0 = +1\langle \phi \rangle_0$ is used. In terms of these fields, the $W^\pm$, $Z^0$ and the photon are given by

$$
W^\pm = \frac{1}{\sqrt{2}}(W_1 \mp iW_2), \quad Z = \frac{gW_3 - g'B}{\sqrt{g^2 + g'^2}} \quad \text{and} \quad A = \frac{g'W_3 + gB}{\sqrt{g^2 + g'^2}},
$$

(1.13)

with masses

$$
M_{W^\pm}^2 = \frac{v^2g^2}{4}, \quad M_Z^2 = \frac{g^2 + g'^2}{4}v^2 \quad \text{and} \quad M_\gamma = 0,
$$

(1.14)

respectively. The weak mixing angle, $\theta_W$, is normally defined as

$$
\tan \theta_W = \frac{g}{g'}
$$

(1.15)
and from this one can write the covariant derivatives in terms of the physical fields as

\[ D_\mu = \partial_\mu + \frac{ig}{\sqrt{2}}(T_+W^+_\mu + T_-W^-_\mu) \]

\[ + \frac{ig}{\cos \theta_W}(T_3 - Q \sin^2 \theta_W Z_\mu) + ieQA_\mu \]

(1.16)

where \( T_\pm = T_1 \pm iT_2 \).

1.6 Non-Minimal Higgs Sectors

The Standard Model uses one scalar doublet to spontaneously break the \( SU(2)_L \otimes U(1)_Y \) symmetry, but there is no physical reason which prevents the Higgs sector from having a richer structure. In particular, there is no physical reason which forbids more than one scalar doublet. In fact, non-minimal Higgs sectors have been used to provide mechanisms for electroweak CP violation [11], strong CP violation [12], and neutrino mass generation [13], and may have a role in \( B^0\bar{B}^0 \) mixing [14]. Also, minimal supersymmetric (SUSY) [15] models require that the Higgs sector be non-minimal.

This thesis treats the minimal extension of the Standard Model, i.e., the Standard Model with the addition of one more complex Higgs doublet field. The results of the previous sections apply here, but now two complex Higgs doublets give eight degrees of freedom and (as before) three of these are "eaten" to provide masses to the \( W^\pm \) and the \( Z^0 \), leaving five remaining degrees of freedom, or five Higgs bosons. These are comprised of two neutral scalars \( H^0 \) and \( h^0 \), one neutral pseudoscalar \( A^0 \), and two charged scalars \( H^+ \) and \( H^- \). In the
search for $H^\pm$ the production cross section was calculated using in the context of electroweak theory. In the search for the pair-produced neutral Higgs, however, one must restrict oneself to more specific models to calculate the cross section; here minimal SUSY is used. Minimal SUSY also gives mass relations between these five Higgs bosons: $M_{H^\pm} > M_{W^\pm}$, $M_{H^0} > M_{Z^0}$, $M_{h^0} < M_{Z^0}$, and $M_{A^0} > M_{h^0}$. Therefore, in this framework the $h^0$ and $A^0$ may be pair-produced while the $H^\pm$ cannot at $\sqrt{s} = M_{Z^0}$.

1.7 Properties of the Higgs Boson

1.7.1 Higgs Production Mechanisms

Higgs bosons at LEP can be pair-produced by coupling to both the intermediate state $\gamma$ or $Z^0$. The Feynman diagrams for these processes are shown in Fig. 1.2. Neutral Higgs pair production is restricted to two of the three possible final states—$h^0A^0$ and $H^0A^0$—since CP-invariance forbids a $Z^0h^0H^0$ vertex. Furthermore, identical Higgs bosons cannot be pair-produced since this would violate Bose symmetry, nor, for models having only Higgs doublets, can the $H^\pm$ be produced singly at tree level via a $Z^0W^\pmH^\mp$ vertex [16].

To lowest order, the cross section for pair-production of the charged Higgs scalar is:

$$\sigma_{H^+H^-} = \frac{\beta^2 \pi \alpha^2}{3 s} \times \left[ 1 - \frac{s(s - M_{Z^0})^2 C_Y C_Y'}{(s - M_{Z^0})^2 + M_{Z^0}^2 \Gamma_Z^2} + \frac{s^2 (C_Y^2 + C_A^2) C_Y'^2}{(s - M_{Z^0})^2 + M_{Z^0}^2 \Gamma_Z^2} \right]$$

(1.17)
Figure 1.2: Pair production of Non-Standard Higgs bosons at LEP.

where

\[ C_V = \frac{1 - 4 \sin^2 \theta_W}{4 \sin \theta_W \cos \theta_W} \quad C_A = \frac{-1}{4 \sin \theta_W \cos \theta_W} \quad C_V' = \frac{-2 + 4 \sin^2 \theta_W}{4 \sin \theta_W \cos \theta_W} \quad (1.18) \]

and \( \beta = \sqrt{1 - 4m^2/s} \). At LEP center-of-mass energies (\( \sqrt{s} = M_{Z^0} \)) the third term of Eq. 1.17 dominates and the cross section at these energies may be approximated as

\[ \sigma_{H^+H^-} \approx \beta^3 \frac{\pi \alpha^2}{3} \left[ \frac{(C_V^2 + C_A^2)C_V'}{\Gamma_2'2} \right]. \quad (1.19) \]

The cross section for pair-production of the neutral Higgs scalar is model-
dependent. The following prediction of minimal SUSY model is used:

\[ \sigma_{h(H^0)A^0} = \frac{1}{4} \left[ \lambda \left( 1, \frac{m_{h(H^0)}^2}{s}, \frac{m_{A^0}^2}{s} \right)^{3/2} \times \frac{4\pi \alpha^2}{3} \left[ \frac{s(C_Y^2 + C_A^2)C_Y^2}{(s - M_{Z^0}^2)^2 + M_{Z^0}^2 \Gamma_L^2} \right] \right] \]  \tag{1.20}

where \( \lambda \) is the phase space function

\[ \lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc \]  \tag{1.21}

and

\[ C_Y' = \frac{-C_H}{\sin \theta_W \cos \theta_W}, \]

where

\[ C_H = \frac{\cos(\alpha - \beta)}{2}, \quad \text{for } e^+e^- \rightarrow H^0 A^0 \]
\[ = \frac{\sin(\alpha - \beta)}{2}, \quad \text{for } e^+e^- \rightarrow h^0 A^0. \]  \tag{1.22}

The mixing angles \( \alpha \) and \( \beta \) are given by

\[ \tan \beta = \frac{v_2}{v_1} \]
\[ \tan 2\alpha = \left( \frac{M_{h^0}^2 + M_{A^0}^2}{M_{A^0}^2 - M_{h^0}^2} \right) \tan 2\beta. \]  \tag{1.23}

where \( v_{1,2} \) are the vacuum expectation values which give masses to the down and up type fermions, respectively.
1.7.2 Higgs Decay Mechanisms

The decay of the Standard Model Higgs boson is distinguished by its characteristic dependence on the mass of its decay products. The decay width of the standard Higgs into fermions is

$$\Gamma(H \rightarrow f \bar{f}) = \frac{N_c g^2 M_H m_f^2}{32\pi M_W^4} \beta_f^2,$$  \hspace{1cm} (1.24)

where $N_c$ is the color factor and $\beta_f = (1 - 4m_f^2 M_H^2)^{1/2}$ is the kinematic factor [18]. For the Standard Model with an additional Higgs doublet, the decay width of the charged Higgs into fermions also depends on the ratio of the vacuum expectation values associated with the two doublets, $\tan \beta = v_2/v_1$ [20]:

$$\Gamma(H^\pm \rightarrow \nu\ell) = \frac{G_F \sqrt{2}}{8\pi} M_{H^\pm} m_\ell^2 \tan^2 \beta$$  \hspace{1cm} (1.25)

$$\Gamma(H^\pm \rightarrow u_{i=1,2}d_{j=1,2,3}) = \frac{G_F \sqrt{2}}{8\pi} M_{H^\pm} 3|V_{ij}|^2 \times \left[ m_i^2 \cot^2 \beta + m_j^2 \tan^2 \beta \right]$$  \hspace{1cm} (1.26)

where $\ell = (e, \mu, \tau)$, $u =$ up-type quarks and $d =$ down-type quarks. From Eq. (1.26), the relative rate $\Gamma(cb)/\Gamma(cs) \simeq 0.01$ to 0.05 for $\tan \beta \simeq 1$, and 0.6 to 2.4 for $\tan \beta \gg 1$, using quark masses and Kobayashi-Maskawa matrix elements as given in Ref. 19. Hence, in the mass range explored in this thesis the following processes will dominate:

$$Z^0 \rightarrow \left\{ \begin{array}{ll}
H^+ H^- \rightarrow \nu \bar{\tau} \bar{\nu} \tau & \text{or} & H^+ H^- \rightarrow \nu \bar{\tau} \bar{c} \bar{b} \\
H^+ H^- \rightarrow c \bar{c} \bar{s} \bar{s} & \text{or} & H^+ H^- \rightarrow c \bar{b} \bar{c} \bar{b}.
\end{array} \right.$$  \hspace{1cm} (1.27)

For the neutral supersymmetric Higgs, this mass dependence is modified by a further dependence on the mixing angles $\alpha$ and $\beta$ defined in Eq. 1.23. Defining
the $H-ff'$ coupling as $c = x + y\gamma_5$, one can write the partial width generally as [18]:

$$\Gamma(H \rightarrow ff') = \frac{N_c}{8\pi M_H} \left[ (x^2 + y^2)(M_H^2 - m_f^2 - m_f^2) - 2(x^2 - y^2)m_f m_f' \right]$$

$$\times \left[ 1 - \frac{2m_f^2}{M_H^2} - \frac{2m_f^2}{M_H^2} - \frac{2m_f^2 m_f^2}{M_H^2} + \frac{m_f^2}{M_H^2} + \frac{m_f^4}{M_H^2} \right]^{1/2}. \quad (1.28)$$

This equation indicates that for neutral Higgs masses greater than twice the mass of the $b$ quark, the dominant decay mode will be $h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}$. In this mass region there is also a contribution from the decay modes $h^0 A^0 \rightarrow b\bar{b} \ \tau\bar{\tau}$ and $h^0 A^0 \rightarrow \tau\bar{\tau} \ \tau\bar{\tau}$ amounting to a small fraction of the total decay width. Only the mode $h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}$ is explored in this thesis.

### 1.7.3 Current Experimental Limits

Charged Higgs with masses below 19.0 GeV have been excluded by the PETRA experiments [23]. There are no previous mass limits on the pair production of the neutral SUSY Higgs $h^0$ and $A^0$, but limits do exist for their production outside of the SUSY framework [24].
Chapter 2

The LEP Storage Ring and the ALEPH Detector

The Large Electron Positron (LEP) storage ring at CERN had its first fill for the pilot run on 13 August 1990. The first $Z^0$ candidate event was observed by the ALEPH detector one day later. On 22 September, the first physics run started, and ALEPH and the other three LEP detectors began collecting $Z^0$ events at the combined rate of roughly 1000 $Z^0$ events/day. In the first section of this chapter, the LEP storage ring is briefly described [25]. A detailed description of the ALEPH detector follows, and some of the detector’s current performance parameters are given [26].

2.1 The LEP Storage Ring

The LEP ring has been designed to accelerate and store counterrotating bunches of electrons and positrons and collide them at up to eight interaction points around the ring. It can ultimately operate at a maximum energy of roughly 125 GeV per beam. In the first (current) phase of operation, “LEP-I,” the
accelerator is running with four bunches and up to 55 GeV per beam but is concentrating at the $Z^0$ peak. Near this peak the luminosity, defined as

$$ L = \frac{N^2 k f}{4\pi \sigma_x \sigma_y}, \quad (2.1) $$

is expected to reach a few $10^{31}$ cm$^{-2}$s$^{-1}$, where $N$ is the number of electrons or positrons per bunch, $k$ is the number of bunches, $f$ is the revolution frequency, and $\sigma_x$ and $\sigma_y$ are the r.m.s. horizontal and vertical beam radii at the interaction point. Only data collected during LEP-I have been analyzed in this thesis; LEP-II data (which will probably be collected at $\sqrt{s} \approx 200$ GeV) will not be available until several years from now.

The ring itself is 26.66 kilometers in circumference and positioned so that it just touches the foot of the French Jura mountains in order to avoid as much contact as possible with solid rock. Also, the entire ring is inclined at a 1.42% slope to further optimize its geological placement.

LEP currently has four bunches per beam, each with a circulation period of 88.9 $\mu$s. One hundred twenty-eight ambient temperature accelerating cavities provide 16 MW of RF power. At the LEP injector LINAC (LIL), positrons are produced in a 200 MeV LINAC via collisions between a beam of electrons and a tungsten converter target. A low-intensity electron gun near this converter produces the electrons. Both types of particles are then sent to a 600 MeV LINAC for further acceleration.

Pulses from the 600 MeV LINAC are stored in the EPA (Electron Positron Accumulator) which converts the repetition rate and bunch charge of the LINACs
to values suitable for the machines downstream. After the EPA, both bunches are sent to the PS (Proton Synchrotron) where they are accelerated to 3.5 GeV, and then they are sent to the SPS (Super Proton Synchrotron) where they attain energies of 20 GeV. Finally, the bunches are sent to the LEP ring where they are accelerated to their full energies. A diagram of the LEP machine is shown in Fig. 2.1.

The LEP lattice consists of 3304 dipole magnets to bend the beams, 816 quadrupole magnets for focusing the beams, and 504 sextupole magnets for chromaticity corrections. All of these magnets operate at ambient temperature. Eight superconducting quadrupoles, two for each of the four current LEP experiments, serve to squeeze the beams just before they collide in order to increase the luminosity. The LEP beams circulate in a vacuum of $10^{-9}$ torr and are expected to reach currents of about 3 mA.

### 2.2 The ALEPH Detector

[The Aleph is] the only place on earth where all places are—seen from every angle, each standing clear, without any confusion or blending.

—from The Aleph, by Jorge Luis Borges
Figure 2.1: The LEP machine: LIL, EPA, PS, SPS, and LEP.
2.2.1 Introduction

The ALEPH detector has been designed to record the result of $e^+e^-$ collisions at center-of-mass energies of $\sqrt{s} \approx M_{Z^0}$, and eventually at center-of-mass energies of about 200 GeV. It is a general purpose detector, built to handle complex high-multiplicity events as well as events with simpler leptonic signatures. It accomplishes these twin objectives by having nearly $4\pi$ solid angle coverage and high resolution tracking detectors coupled with finely granulated calorimetry.

The ALEPH detector is located in Echenevex, France, 143 meters underground at the foot of the Jura mountains. It is approximately 16 meters high and 21.4 meters wide, and it weighs about 3000 metric tons. It is comprised of many subdetectors, most of which are arranged as concentric cylinders with axes parallel to the beam axis. The subdetectors of the apparatus, increasing radially, are the VDET (Vertex DETector; not yet fully installed), ITC (Inner Tracking Chamber), the TPC (Time Projection Chamber), the ECAL (Electromagnetic CALorimeter), the HCAL (Hadron CALorimeter and return yoke for the magnetic field), and the muon chambers. The LCAL (Luminosity CALorimeter) and SATR (Small Angle TRacking device) are located in the regions near the beam pipe at both ends of the detector. A cut-away view of the apparatus is shown in Fig. 2.2, and in the following paragraphs the individual subdetectors are described in detail.
Figure 2.2: The ALEPH apparatus and its subdetectors: luminosity monitor (1), inner tracking chamber (2), time projection chamber (3), electromagnetic calorimeter (4), superconducting coil (5), hadronic calorimeter (6), and muon chambers (7). Also shown are LEP's low-beta superconducting quadrupoles (8).
2.2.2 The ITC

The inner tracking chamber is a two meter long cylindrical multiwire drift chamber operating at atmospheric pressure in Ar(50%) + ethane(50%) gas. Small quantities of alcohol and water are added to inhibit formation of polymers and deposition of dirt on the wires. At an operating voltage of 2.1 kV the gas gain is about $2 \cdot 10^4$. The ITC provides up to eight accurate $r$-$\phi$ points between 16-24 cm in radius for tracking, and it is also used in the Level-I trigger since its readout is very fast (2-3 $\mu$s). The detector consists of 960 sense wires and 3840 field wires parallel to the beam direction and arranged in hexagonal drift cells. These cells are small (4.7-6.5 mm drift distances) to permit the fast readout needed for the trigger. The wire configuration is shown in Fig. 2.3. The $r$-$\phi$ coordinate is measured using the drift time, and the $z$ coordinate is measured using the difference in the arrival times of pulses at the ends of each sense wire. The resolution is 150 $\mu$m in $r$-$\phi$ and 3 cm in $z$.

2.2.3 The TPC

As the main tracking chamber, the time projection chamber is in many ways the heart of the ALEPH detector. It gives excellent 3-dimensional track separation in the high-multiplicity environment of LEP, high momentum resolution and very good particle identification with its $dE/dz$ measurement.

The chamber is a cylindrical device 4.4 meters long and 3.6 meters in diameter whose axis lies parallel to the beam axis. It is operated at atmospheric pressure in 9% Ar–91% CH$_4$ gas. The drift field is generated by a membrane at the center
Figure 2.3: A section of the ITC, showing its wires arranged in hexagonal drift cells.

of the TPC which divides the detector into two identical drift volumes. This
field is kept uniform by graded potentials provided by conducting strips on the
inner and outer field cages. With the solenoidal magnetic field, the electric and
magnetic fields are very nearly parallel (or antiparallel) throughout the TPC.

Charged particles passing through the TPC ionize the gas within, leaving
behind a trail of electrons. These electrons drift along the electric field lines of
the drift field and the magnetic field lines of the solenoid to one of the two end-
plates, where they are collected. Individual proportional wire chambers called "sectors," of which there are 18 per endplate, measure the time of arrival and position of these electrons, yielding a 3-dimensional space point, and the amount of ionization deposited, yielding a measurement of the $\frac{dE}{dz}$ . Fig. 2.4 shows a cut-away view of the TPC.

Figure 2.4: The ALEPH TPC, showing the central high-voltage membrane, the field cages, the endplates and their individual sectors, and the solenoid (the ECAL is not shown for clarity).

At each sector the spatial and $\frac{dE}{dz}$ data are developed in the following manner: the drifting electrons pass through two wire grids and are then proportionally amplified by the high voltage sense wires at the third and final wire grid.
The cathode plane, situated directly beneath the sense wire grid, consists of precisely located $6 \times 30$ mm$^2$ metallic pads arranged in 21 circular rows concentric with the interaction point. (Between these pad rows are special pads measuring $6 \text{ mm} \times 15^\circ \text{ in } r \times \phi$ which are used in the Level-II trigger.) Avalanches on the sense wires induce signals on these pads. The TPC reads out a total of about 41,000 pads and 6,400 wires. An end-on view of an endplate with its 18 sectors is shown in Fig. 2.5. The arrangement of the pad plane and the wire grids in each sector is shown in Fig. 2.6.

Figure 2.5: One of the two TPC endplates. The 6 inner sectors are staggered with respect to the 12 outer sectors, and the latter dovetailed with respect to one another, in order to reduce the likelihood of a charged particle escaping detection.
Figure 2.6: A TPC sector edge, showing the pad plane and the three wire grids. The wires are glued to G10 substrate at the sector edges and soldered to printed-circuit traces through which they are connected to the external world. In between the grids are conducting strips used to compensate for spatial and gain distortions near the sector edges.

From the drift time and velocity of the ionized electrons, the $z$-coordinate at the point where the ionization is produced is calculated, and from an interpolation of the signals induced on the pads its $r$-$\phi$ position is obtained. The radial position is simply the radius of the pad row whose pads received signals. A maximum of 21 3-dimensional coordinates and a maximum of about 350 $dE/dz$ samplings can be measured.

A track in the TPC describes a helix whose projection onto the endplate is a
circular arc. By measuring the sagitta of this arc one can calculate the radius of curvature of the particle in the known magnetic field and hence its momentum and charge. The momentum resolution is

\[ \frac{\delta p_t}{p_t} = 0.0015 \cdot p_t. \]  

(2.2)

This resolution improves to \(0.0010 \cdot p_t\) when the ITC measurement is combined with that of the TPC, as shown in Fig. 2.7.

![Graph showing momentum resolution](image)

Figure 2.7: The momentum resolution of the TPC as measured by a data sample of highly collinear dimuon events. The Gaussian fits give \(\sigma = 0.0487 \pm 0.0013\) and \(\sigma = 0.0483 \pm 0.0006\) for \(Q = -1\) and \(Q = +1\), respectively.

The other two wire grids mentioned above are called the ground (or cathode) grid and the gating grid. The ground grid is the middle grid and its function is to
terminate the drift field and isolate it from the high fields developed at the sense wires. The gating grid is the wire grid closest to the active volume of the TPC, and its function is to prevent the positive ions created at the sense wires during avalanches from entering the active volume. This is very important for the integrity of the operation of the TPC since these ions can form localized clouds of positive charge which can, in turn, distort the drift field and significantly degrade the TPC momentum resolution. The principle, implementation and performance of the gating system is described in Appendix A.

At the edge of each TPC sector, where the ends of the sense wires are in close proximity to the edge's dielectric G10 substrate, the electric field the wires generate is no longer uniform. This causes non-uniformity in the gain at the sector edge and might cause spatial distortions as well if the electrons do not drift mainly along the magnetic field lines in this region. Anticipating this effect, the TPC sectors were built with three correction strips (see Fig. 2.6) on which three independent DC voltages can be placed. Measurements have shown that one can reduce the gain distortion at the sector edge to negligible values by placing the appropriate voltages on these strips (see Appendix B for details).

2.2.4 The ECAL

The electromagnetic calorimeter is a very finely-grained lead/wire chamber sampling device. It consists of a 12-module barrel and two 12-module endcaps which, taken together, completely surround the TPC and are themselves completely contained by the solenoid (see Fig. 2.8). Each module has 45 lead/wire chamber
layers, giving a total nominal thickness of 22 radiation lengths. Each layer consists of anode wires in 80% Xe + CO₂ gas and 30 × 30 mm² pads. A schematic of such a layer is shown in Fig. 2.9. The anode wires amplify ionization resulting from showers developed in the lead sheets, and these amplified signals then couple to the nearby pads. The pads are internally arranged such that they form geometrical projective towers subtending 1° × 1° sin θ in solid angle in the barrel. The signals from a plane of wires are ganged together and read out as a single channel, giving a total of 1620 wire channels. These wire signals are used in the trigger and as a check on the energy measured by the pads. Each of the
Figure 2.9: A layer of the ECAL, showing the anode wires, pads, and lead sheet.

~ 70,000 towers is split into three storeys in depth for readout (corresponding to the first four, middle nine, and last nine radiation lengths), giving a total of more than 200,000 pad channels.

As measured by test beam data, the ECAL achieves \((16.5 \pm 0.3)\%/\sqrt{E}\) resolution, with a constant term of \((0.3 \pm 0.1)\%\) for electromagnetically-induced showers. Its angular resolution in \(\theta\) and \(\phi\) is \((17.6 \pm 0.9)\%/\sqrt{E}\), with a constant term of \((2.0 \pm 0.2)\%\), measured in tower units. Both these measurements are averages over entry angles ranging from 45°–90°. The hadron misidentification probability, \(\epsilon_r\), for selection criteria which have an electron identification
efficiency of $80 \pm 2\%$, is [27]

$$\epsilon_r < 7 \times 10^{-3}.$$ (2.3)

2.2.5 The Superconducting Solenoid

The solenoidal magnet consists of a superconducting coil which produces a uniform 1.5 Tesla field parallel to the beam direction, and an iron return yoke which serves both as the hadron calorimeter (see next section) and the main support structure of the entire apparatus. The solenoid is designed to provide the homogeneous magnetic field required by the TPC, with its radial thickness kept to a minimum (1.6 radiation lengths) to reduce multiple scattering. This homogeneity requirement may be expressed as:

$$\int_0^{2.2m} B_r/B_z \, dz < 2 \text{ mm}$$ (2.4)

$$r < r_{TPC}.$$

Large values of this integral correspond to distortions in the measured sagitta (weighted by $1/\omega \tau \approx 8$). Measurements have shown that this integral is never greater than about 1.5 mm. The same measurements also included a precise mapping of the entire field in the TPC volume, with a resolution of $\delta B/B \approx 10^{-4}$, which is used to find the correct current to place in the magnet’s compensating coils. The mapping is also used to correct the data for the small remaining field inhomogeneities.
2.2.6 The HCAL and the Muon Chambers

The HCAL consists of 22 iron slabs, each 5 cm thick, and an additional outer slab 10 cm thick, comprising a total thickness of 1.2 m or 7.2 interaction lengths. Between these slabs are gaps containing limited streamer mode tubes, for a total of 22 layers of tubes. The tubes operate at 4250 V in a gas consisting of Ar::CO₂::isobutane in the ratio 13::57::30. A schematic diagram of the HCAL is shown in Fig. 2.10. Hadronic showers are developed in the iron and the resulting particles get detected by the tubes. The avalanches induce signals on nearby pads and on conducting strips running parallel to the tube wires. Like the ECAL, the pads are arranged in a projective geometry, but each tower subtends a larger solid angle in the barrel: 4° × 4° sin θ. The resolution of the HCAL is roughly 85%/√E. The signals from the wires themselves are used in the trigger. The conducting strips provide a digital pattern corresponding to particle penetration depth which is used for muon identification. An enlargement of a section of an HCAL module, showing the pads, the tubes with their sense wires, and the conducting strips is shown in Fig. 2.11.

Two additional double layers of streamer tubes, one just outside the HCAL, the other 50 cm beyond, comprise the muon chambers, which can measure the exit angle of muons to an accuracy of σ = 10 mrad. However, at the present time only one double layer is installed. A schematic of one of these double layers, showing the tubes and the z and y readout strips, is shown in Fig. 2.12.
Figure 2.10: The ALEPH hadron calorimeter.
Figure 2.11: An enlargement of a module of the ALEPH hadron calorimeter, showing the pads, the tubes with their sense wires, and the conducting strips.
Figure 2.12: A schematic of one of the double layers of the muon chambers, showing the tubes and the $x$ and $y$ readout strips.
2.2.7 The Luminosity Monitors

The luminosity measurement is obtained primarily from the LCAL, which is very similar in construction and readout to the ECAL. It is located ±2.7 meters on either side of the interaction point, between the beam pipe and the ECAL endcap. The device covers an angular range between 55 and 155 mrad. Four semicylindrical modules, two on each side of the interaction point, contain 38 lead/wire chamber layers, for a total radiation length of 24.6 radiation lengths. Its energy and position resolutions are 20%/√E and 2.5 mm, respectively.

The SATR has an angular resolution of 0.05 mrad and is located ±2.45 meters on either side of the interaction point. The device covers an angular range between 40 and 90 mrad. Thirty-six half-planes of drift tube chambers are arranged in four sections (two on each side of the interaction point), and each half-plane has four 45° sectors, each of which have fourteen drift tubes. Four groups of nine half-planes are rotated 15° with respect to each other to avoid dead zones.

With both the LCAL and the SATR, a 1.3% systematic error is achieved on the absolute luminosity. This error is low by virtue of the excellent measurement of the energies of the $e^+e^-$ Bhabha pairs and of the shower position.
Chapter 3

Event Triggering, Reconstruction, and Simulation

A detector of ALEPH's complexity requires a sophisticated and flexible triggering system to keep the size of the data sample reasonably small while simultaneously enhancing its physics content. All data satisfying the trigger requirements are analyzed by a general reconstruction program, in which the data from each subdetector are reconstructed separately and then relationships between the data from the subdetectors are made (e.g., an attempt is made to join the signals left by a charged particle in the ITC and TPC into a single track). The triggering system and the reconstruction code are described in the first section of this chapter. The other section of this chapter is concerned with the simulation of the detector and the relevant physics processes, which are essential to the understanding of the detector itself and to the behavior of background processes and the signals of interesting physics processes.
3.1 Event Triggering and Reconstruction

The data acquisition system must strike a balance between the desire to read out any event with potentially interesting physics content and the speed and storage limitations of the detector electronics and associated computers. To reduce the size of the data sample while enriching its physics content, ALEPH uses a triggering system split into two levels [26]. The “Level-I” trigger decides whether or not to initiate event digitization within 5\(\mu s\) after the beam crossing. It consists of up to 32 different triggers, but only the following are relevant for triggering on multihadronic \(Z^0\) decays:

1. an energy deposit of at least 6.5 GeV in the barrel of the ECAL;
2. an energy deposit of at least 3.8 GeV in one of the two endcaps;
3. an energy deposit of at least 1.6 GeV in each of the endcaps;
4. an ITC track (defined as a signal in at least 5 of the 8 cylindrical wire layers recording ionization from charged particles) in azimuthal coincidence with a deposit of at least 1.3 GeV in the ECAL;
5. an ITC track in azimuthal coincidence with a signal penetrating at least 40 cm of iron in the HCAL.

Since these triggers are independent, and since events are often accepted on the basis of more than one of them, individual trigger efficiencies can be readily
measured. These efficiencies are $99.96 \pm 0.02\%$ for the three ECAL energy triggers [38], $96.84 \pm 0.42\%$ for the ITC-ECAL trigger, and $90.76 \pm 0.92\%$ for the ITC-HCAL trigger [39] for hadronic and leptonic $Z^0$ decays.

To measure the luminosity from low-angle Bhabha events, used to determine the expected number of Higgs signal events in the data sample, three additional conditions can activate the Level-I trigger:

1. a coincidence of 20 GeV deposited on one side with 16 GeV on the other side of the LCAL, with no azimuthal correlation requirement;

2. a deposit of at least 31 GeV on either side of the LCAL ("single-arm" triggers);

3. pre-scaled single-arm triggers with both 16 and 20 GeV thresholds to provide an estimate of the beam-related backgrounds.

Trigger efficiencies are $99.7 \pm 0.2\%$ for Bhabha events.

The "Level-II" trigger uses TPC information, available roughly 45$\mu$s after beam crossing, to determine whether or not at least one track comes from the origin. In typical operating conditions, 20% of Level-I triggers are track-only triggers, and of these about 75% are rejected by Level-II. Rejected events are normally come from beam-gas or cosmic ray backgrounds. The Level-II trigger was used only at the end of the 1989 run to veto ITC-ECAL and ITC-HCAL triggers.
In spite of the large improvement in the signal-to-noise ratio provided by the Level-I and Level-II triggers, many unwanted events survive and must be removed by offline selection criteria. These criteria are applied after all surviving events are processed through JULIA, the reconstruction program. JULIA joins coordinates provided by the tracking detectors to reconstruct the tracks left by charged particles, and clusters the contiguous deposits in the calorimetry to form calorimeter "objects." Those objects which are close to the extrapolated position of a reconstructed track(s) are associated to them for use in offline analyses, and isolated objects are kept as well. The criteria which select multihadronic $Z^0$ decays and remove unwanted events are discussed in Chap. 4.

3.2 Event Simulation

The Higgs signals searched for in this thesis are multi-jet processes which deposit a substantial amount of energy in the detector. For this reason, the only relevant background comes from $Z^0 \rightarrow q\bar{q}(g)$ processes. All of these processes are passed through the full simulation program of the ALEPH detector, "GALEPH," which includes effects of particle interactions with elements of the detector itself, trigger and detection efficiencies, detector geometry, and energy and tracking resolutions, and then reconstructed with the same program used to reconstruct the data, "JULIA." For all processes initial state radiative corrections are made according to the Berends and Kleiss prescription [40].
3.2.1 $Z^0 \rightarrow q\bar{q}(g)$ Background Simulation

The simulation of the $Z^0 \rightarrow q\bar{q}(g)$ background is a two-step process. In the first step a configuration of initial state partons is made using the QCD generator BREM5 [41]. BREM5 includes final state radiative corrections and provides the final state $q\bar{q}(g)(\gamma)$ according to branching ratios predicted by the Standard Model. Feynman diagrams for these processes are shown in Fig. 3.1. In the second step of the simulation, the final state partons generated by BREM5 are fragmented into hadrons. Unstable hadrons resulting from this fragmentation are decayed according to recently published data [19] with modifications to the default branching ratios for D and D* mesons [19]. Routines from LUND JET-SET version 6.3 [42] are used to perform the fragmentation and to decay the unstable hadrons.

Since it is not yet known how to evolve quarks to hadrons, a model for this fragmentation process is needed. In the work presented here, the model known as "parton shower" is used in conjunction with the LUND symmetric [43] and Peterson et. al. fragmentation functions [45]. The parton shower model is based on the Altarelli–Parisi evolution Eqs. [46] which are used iteratively to determine when a branching $q \rightarrow qg, g \rightarrow gg, \text{ or } g \rightarrow q\bar{q}$ occurs. Once the invariant mass of the resulting two–parton systems falls below a specified cutoff, the branching process is terminated and the partons are fragmented.

The partons carry color charge, and as such are subject to the strong confinement force. Consequently, as adjacent partons move apart the potential of
Figure 3.1: Feynman diagrams for $Z^0 \rightarrow q\bar{q}(g)$ processes. Examples are given for a.) standard QCD processes, b.) QCD processes with first order and c.) second order corrections in $\alpha_s$. All these processes can produce multijet signatures in the detector.
the system increases until eventually a new pair of partons are created out of the vacuum. Pictorially, a color flux tube is stretched between the adjacent partons. This flux tube is ascribed an energy per unit length $\kappa \simeq 1 \text{ GeV/fm}$, and typically a break occurs when the tube is 2–5 fm in length.

At each break, the newly created partons are given a transverse momentum following Gaussian distributions in $p_x$ and $p_y$, each of width 0.325 GeV, and the mass of the resulting hadron(s) is determined from flavor and spin considerations. Finally, the longitudinal fragmentation of the hadron(s) is made to determine how large a fraction of the available energy a newly-created hadron takes. This is done for each $q\bar{q}$ system according to a probability distribution $f(z)$, where

$$z = \frac{(E + p_{||})_{\text{hadron}}}{(E + p_{||})_{\text{parton}}}.$$  \hspace{1cm} (3.1)

Here $p_{||}$ is the component of momentum along the original parton direction. A value of $z$, $z_1$, is chosen from $f(z)$ for the first hadron, leaving the fraction $1 - z_1$ for the remainder system. Hadrons are produced until the remainder system reaches some minimum value of $E + p_{||}$. The resulting hadrons form groups of particles called jets since each group moves out in the direction of its original parent parton, but each particle within this group has some transverse momentum, giving them jointly the appearance of a jet of particles in the detector.
Two functional forms for $f(z)$ are used in this simulation, one for the lighter quarks $u, d, s$ and one for the heavier quarks $b, c$. The lighter quarks are fragmented with the LUND symmetric fragmentation function [43]

$$f(z) = z^{-1}(1 - z)^a \exp \frac{-b m_t^2}{z},$$  \hspace{1cm} (3.2)

where $m_t$ is the transverse mass of the hadron $E^2 - p_t^2$. Fits to experimental data give $a = 0.45$ and $b = 0.7$ [44]. The heavier quarks are fragmented with the harder Peterson fragmentation function [45]

$$f_{b,c}(z) = \left[ z \left( 1 - \frac{1}{z} - \frac{\epsilon_{b,c}}{1 - z} \right)^2 \right]^{-1}$$  \hspace{1cm} (3.3)

where $\epsilon_c = 0.025$ and $\epsilon_b = 0.015$.

### 3.2.2 Charged and Neutral Higgs Signal Simulation

The simulations of the processes

$$Z^0 \rightarrow \begin{cases} \text{H}^+ \text{H}^- & \rightarrow \bar{c} \bar{s} \\ \text{H}^+ \text{H}^- & \rightarrow \bar{c} b \\ \text{A}^0 \text{h}^0 & \rightarrow b \bar{b} \end{cases}$$  \hspace{1cm} (3.4)

shown in Fig. 3.2, begin with the selection of a production angle and two decay angles. The production angle is the angle the Higgs pair makes with the original $e^+ e^-$ direction, and it is randomly generated according to a $\sin^2 \theta$ probability distribution, characteristic of the pair production of scalar particles. The decay angles are the angles each of the Higgs' daughter quarks make with the original parent Higgs' direction, and they are randomly generated according to an isotropic probability distribution in the rest frame of the parent Higgs, characteristic of the decay of a scalar particle. Once the direction of the daughter
quarks has been determined, they are boosted back into the $Z^0$ rest frame and given to LUND to fragment as described in section 3.2.1 above.

\[ \begin{aligned}
\text{e}^+ & \quad \text{Z}^0, \gamma \quad \text{e}^- \\
& \quad H^+ \quad \text{c} \\
& \quad H^- \quad \text{\overline{c}} \\
\text{e}^+ & \quad \text{Z}^0 \quad \text{h}^0 \\
& \quad \text{b} \\
& \quad \text{A}^0 \quad \text{b}
\end{aligned} \]

Figure 3.2: Feynman diagrams for the processes listed in Eqn. 3.4. In the upper diagram, $s$ is replaced by $b$ for the mode $H^+ H^- \rightarrow b\overline{b} \quad c\overline{b}$. 
Chapter 4

The $H^+H^-$ and $A^0 h^0$ Four-Jet Analyses

The processes searched for in this thesis,

$$Z^0 \rightarrow \left\{ \begin{array}{l}
H^+ H^- \rightarrow c\bar{s} \ c\bar{s} , \\
H^+ H^- \rightarrow c\bar{b} \ c\bar{b} \ \text{and} \\
h^0 A^0 \rightarrow b\bar{b} \ b\bar{b} ,
\end{array} \right.$$  \hspace{1cm} (4.1)

each would leave a multi-jet signature in the ALEPH detector. In the simplest case, each quark in the final state fragments into its own distinct jet, but often the quarks radiate gluons and fragment in such a way as to render the final state more complex. Furthermore, since only the charged particles are used in these analyses, a considerable amount of information about each jet is lost, adding to the complexity. To overcome these difficulties, a clustering algorithm is used to group the charged particles of multihadronic events into four and only four jets. Then, a procedure by which the four jets are forced to conserve energy and momentum is applied in order to compensate for unused neutral energy. Finally, various selection criteria are applied to enhance the simulated Higgs signals relative to the background, which consists almost exclusively of
$Z^0 \rightarrow q\bar{q}(g)$ events in which the final state multi-jet signature (also clustered into four jets) happens to have a topology mimicking that of one or more of the processes listed above.

The searches for the charged Higgs and neutral SUSY Higgs use similar analyses to enhance the four-jet-like signals characteristic of the processes listed in Eq. 4.1. In both analyses only multihadronic events, $Z^0 \rightarrow \text{hadrons}$, are taken by requiring

$$\text{number of good tracks } \geq 5$$

and

$$\text{total energy of good tracks } \geq 0.1 \cdot E_{\text{cm}},$$

(4.2)

where $E_{\text{cm}}$ is the center-of-mass energy and a good track is defined as one making an angle $\alpha$ with the beam axis such that $|\cos \alpha| < 0.95$, having at least 4 TPC coordinates, $|d_0| < 2 \text{ cm}$, and $|z_0| < 10 \text{ cm}$. Here, $|d_0|$ is defined as the distance of closest approach to the interaction point in the plane perpendicular to the beam axis and $z_0$ is defined as the coordinate along the beam with respect to the interaction point. Only charged particles meeting these criteria are used in these analyses. (Neutral particles measured by the calorimetry are not used because during the 1989 run the ECAL and HCAL were often noisy and/or unreliable.)

With all the good charged particles the aplanarity is calculated. The aplanarity is defined as in Ref. 34, and gives a measure of the flatness of the event. To remove a large fraction of the two- and three-jet background, events with

$$\text{aplanarity } < 0.02$$

(4.3)
are discarded. The aplanarity distributions for the data, the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo, and a simulated 30 GeV charged Higgs are shown in Fig. 4.1. In this figure, as in all figures which follow, both the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo and the simulated charged Higgs signal are normalized to the luminosity of the data (1.17pb$^{-1}$).

![Aplanarity distributions](image)

**Figure 4.1:** The aplanarity distributions for the data, the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo, and a simulated 30 GeV charged Higgs, for events passing the multihadronic event selection criterion.

All the good charged particles in each multihadronic event are then clustered into at least two jets [48]. Events having less than four jets or more than five jets are re-clustered [49] so that they have exactly four jets. Five-jet events are
made into four-jet events by merging the jet having the smallest energy with the jet closest to it spatially. This procedure provides a slightly better resolution on the reconstructed Higgs masses than re-clustering alone since at larger Higgs masses gluon radiation is a more frequent occurrence.

To insure that all four jets are well-reconstructed, each is required to consist of at least three charged particles, and the magnitude of the sum of the charges of these charged particles must not exceed two. An event with any jet failing to meet these criteria is discarded.

The resolution on the reconstructed Higgs mass is improved by a factor of 2-3 and the reconstructed mass itself is more accurately obtained by using “velocity reconstruction,” a procedure which imposes energy and momentum conservation on the four jets in an effort to compensate for unused neutral particles. In velocity reconstruction, the observed jet velocities $\tilde{\beta}_{\text{jet}}$, calculated by assigning all charged particles the mass of the charged pion, are assumed to be equal to the true jet velocities. Then, the energies $E_{\text{jet}}$ of the individual jets are recalculated [35] by solving the equations of energy and momentum conservation

$$\sum_{\text{jet}=1}^{4} \tilde{E}_{\text{jet}} = E_{\text{cm}} \quad \text{and} \quad \sum_{\text{jet}=1}^{4} \tilde{E}_{\text{jet}} \tilde{\beta}_{\text{jet}} = 0. \quad (4.4)$$

Events with any jet having a negative recalculated energy are removed. Also, events with any jet having a recalculated energy $\tilde{E}_{\text{jet}}$ differing greatly from its original observed value $E_{\text{jet}}$, as defined by

$$0.01 \ \text{GeV} > \frac{\tilde{E}_{\text{jet}}}{E_{\text{jet}}} |\tilde{E}_{\text{jet}} - E_{\text{jet}}| > 100.0 \ \text{GeV}, \quad (4.5)$$
<table>
<thead>
<tr>
<th>Selection Criterion</th>
<th>Cumulative Efficiencies (%)</th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Data</td>
<td>$Z^0 \to q\bar{q}(g)$ M.C.</td>
<td>H$^+$ H$^-$</td>
<td>h$^0$ A$^0$</td>
</tr>
<tr>
<td>Multihadronic</td>
<td>97.5</td>
<td>97.5±0.8</td>
<td>100.0</td>
<td>100.0</td>
</tr>
<tr>
<td>4-Jet Clustering</td>
<td>93.5</td>
<td>94.6±2.0</td>
<td>100.0</td>
<td>100.0</td>
</tr>
<tr>
<td>Aplanarity</td>
<td>16.4</td>
<td>16.8±0.1</td>
<td>70.6±1.9</td>
<td>86.9±2.2</td>
</tr>
<tr>
<td>Vel. Rec. Eq. 4.4</td>
<td>11.3</td>
<td>12.6±0.1</td>
<td>59.5±1.7</td>
<td>72.8±2.0</td>
</tr>
<tr>
<td>Vel. Rec. Eq. 4.5</td>
<td>7.8</td>
<td>8.8±0.1</td>
<td>47.0±1.6</td>
<td>60.1±1.8</td>
</tr>
<tr>
<td>Jet Multiplicity</td>
<td>6.2</td>
<td>7.3±0.1</td>
<td>40.9±1.4</td>
<td>53.0±1.7</td>
</tr>
<tr>
<td>Jet Charge</td>
<td>4.5</td>
<td>5.7±0.1</td>
<td>33.9±1.3</td>
<td>42.8±1.5</td>
</tr>
</tbody>
</table>

Table 4.1: The efficiencies for the selection criteria up to and including jet charge, after all previous criteria have been applied, for the data, the $Z^0 \to q\bar{q}(g)$ Monte Carlo, a 30 GeV H$^\pm$ (where H$^+ H^- \to c\bar{c} \bar{s}s$), and a 30, 31 GeV h$^0$ A$^0$ pair.

are removed. Events removed for either of these reasons are typically events with an abnormally large amount of missing energy, or events for which one or more of the jets have a very low energy ($\approx 3$ GeV). For the remainder of each of the four-jet analyses, the recalculated jet energies are used.

The cumulative efficiencies for each of these selection criteria are given in Table 4.1 for the data, the $Z^0 \to q\bar{q}(g)$ Monte Carlo, and simulated 30 GeV H$^\pm$ and 30, 31 GeV h$^0$ A$^0$ Higgs signals.

The charged Higgs and SUSY neutral Higgs analyses require many different selection criteria in addition to those mentioned above due to the copious $Z^0 \to q\bar{q}(g)$ background. As mentioned earlier, these selection criteria are very
similar since both processes are multihadronic, 4-jet-like processes. However, there are some crucial differences:

1. The cross sections differ significantly, with that for $Z^0 \rightarrow H^+ H^-$ less than that for $Z^0 \rightarrow h^0 A^0$ for similar masses. The expected number of events in the 1989 data sample as a function of Higgs masses for $H^+ H^-$ and $A^0 h^0$ are shown in Figs. 4.2 and 4.3, respectively. In these calculations, initial state radiative corrections have been included as discussed in section 3.2. Figs. 4.4 and 4.5 show the branching fractions $\Gamma(Z^0 \rightarrow H^+ H^-)/\Gamma_Z^{tot}$ and $\Gamma(Z^0 \rightarrow h^0 A^0)/\Gamma_Z^{tot}$, respectively.

2. The charged Higgs satisfy the important constraint $M_{H^+} = M_{H^-}$, while $M_{A^0}$ and $M_{h^0}$ can differ substantially. This constraint improves the resolution on the reconstructed mass and reduces combinatorical background in the charged Higgs analysis.

3. In the mass range explored in this thesis, the dominant decay mode for the SUSY neutral Higgs is

$$h^0 A^0 \rightarrow b\bar{b} b\bar{b}. \quad (4.6)$$

With 4 $b$ quarks in the final state, a significant number of high $p_T$ leptons from semi-leptonic $b$ decays are expected (here, $p_T$ is defined as the momentum of the lepton transverse to its nearest jet). In contrast, the charged
Higgs decay modes are

\[ H^+ H^- \rightarrow c \bar{s} \bar{s} \]
\[ H^+ H^- \rightarrow c \bar{s} \bar{b} \]
\[ H^+ H^- \rightarrow c \bar{b} \bar{b} \]  \hspace{1cm} (4.7)

(plus conjugates), and so fewer high \( p, p_t \) leptons from semi-leptonic decays are produced. By using electron and muon identification in the SUSY neutral Higgs analysis, a large gain in signal-to-noise ratio \( (S/N) \) is obtained, and so the selection criteria for this four-jet-like process are generally be less restrictive than those of the otherwise similar charged Higgs analysis.
Figure 4.2: The expected number of $Z^0 \rightarrow H^+ H^-$ decays in the 1989 data sample as a function of $M_{H^\pm}$. 
Figure 4.3: The expected number of $Z^0 \rightarrow h^0 A^0$ decays in the 1989 data sample as a function of $M_{A^0}$ for selected values of $M_{h^0}$ (recall that $M_{A^0} > M_{h^0}$, and that the cross section decreases both with increasing $M_{A^0}, M_{h^0}$ and increasing $M_{A^0} - M_{h^0}$).
Figure 4.4: The branching fraction $\Gamma(Z^0 \to H^+H^-)/\Gamma_Z^{tot}$ as a function of $M_{H^\pm}$. 
Figure 4.5: The branching fraction $\Gamma(Z^0 \to h^0 A^0) / \Gamma^\text{tot}$ as a function of $M_{A^0}$, for several values of $M_{h^0}$. 
4.1 Charged Higgs Selection Criteria

With four jets in the final state there are three possible jet-jet pairings. Labelling the jets \(i, j, k\) and \(l\), these pairings are

\[
i j \ kl, \ ik \ jl, \ and \ il \ jk. \quad (4.8)
\]

Since the charged Higgs satisfy \(M_{H^+} = M_{H^-}\), they should each have half the center-of-mass energy \(E_{cm}/2\). Hence, only the jet-jet pairing minimizing

\[
\left| \vec{E}_i + \vec{E}_j - \frac{E_{cm}}{2} \right| \quad (4.9)
\]

is taken, where \(i, j\) denote two of the four available jets. From a simulation of \(H^+ H^- \rightarrow c \bar{s} \ \bar{c} s\) with \(M_{H^\pm} = 30\ GeV\), it is estimated that the true combination is correctly chosen in \((73.3 \pm 2.3)\%\) of the events. For the chosen pairing, with \(i, j, k, l\) denoting the individual jets and \(ij\) and \(kl\) denoting the two jet-pair systems, the following quantities are then calculated:

1. the invariant masses \(M_{ij}\) and \(M_{kl}\) of the two jet-pair systems;
2. the production angle \(\theta^p\), defined as the acute angle between the jet-pair system directions and the beam axis;
3. the two opening angles \(\theta_{ij}^{open}\) and \(\theta_{kl}^{open}\), defined as the angles made between the jets in each jet-pair system in the \(Z^0\) rest frame;
4. the two decay angles \(\theta_{ij}^d\) \(\left(\theta_{kl}^d\right)\), defined, in the \(ij(kl)\) rest frame, as the angles the boosted jets \(i, j(k, l)\) make with the \(ij(kl)\) direction in the \(Z^0\) rest frame;
5. the boosted sphericities $S_{ij}^b(S_{kl}^b)$, defined as the sphericities in the $ij(kl)$ rest frame (the sphericity itself is defined as in Ref. 34).

The production angle, the opening angles, and the decay angles are illustrated in Fig. 4.6.

![Diagram](image)

Figure 4.6: The production angle $\theta^p$, the opening angle $\theta^\text{open}_{ij}$, and the decay angle $\theta^d_{ij}$ for jets $i, j, k$ and $l$.

The two charged Higgs particles have the same mass, and therefore the requirement $|M_{ij} - M_{kl}| < 4 \text{ GeV} \approx 2\sigma \overline{M}$ is imposed, where $\sigma \overline{M}$ is the resolution for the average reconstructed mass, $\overline{M} = \frac{1}{2}(M_{ij} + M_{kl})$. The quantities $|M_{ij} - M_{kl}|$ and $\overline{M}$ are shown in Figs. 4.7 and 4.8, respectively. Fig. 4.8 shows
that the resolution on the average reconstructed mass is $\sigma_{\bar{M}} \approx 2.0$ GeV. In the plots which follow (except Figs. 4.8 and 4.14) all events are required to have $\bar{M} = 30.0 \pm 2.5$ GeV (i.e., $27.5 \leq \bar{M} \leq 32.5$ GeV), making the comparisons which are drawn between the data, the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo background, and a simulated 30 GeV charged Higgs signal with $H^+ H^- \rightarrow c\bar{s} \bar{c}s$ consonant with the procedure used to set the limit in Chapter 5.

Figure 4.7: The mass difference $|M_{ij} - M_{kl}|$ for multihadronic events satisfying aplanarity $> 0.01$ and $\bar{M} = 30.0 \pm 2.5$ GeV, and minimizing Eq. 4.9. Events are accepted if they satisfy $|M_{ij} - M_{kl}| < 4$ GeV.
Figure 4.8: The average reconstructed mass $\bar{M}$ for multihadronic events satisfying aplanarity $> 0.01$ and minimizing Eq. 4.9. The resolution $\sigma_{\bar{M}}$ using a Gaussian fit to the region between 25–35 GeV is $2.0 \pm 0.06$ GeV and the mean of the fit is $29.2 \pm 0.03$ GeV. The mean is low by 0.8 GeV due to the removal of tracks which do not satisfy the criteria for good tracks discussed at the beginning of this chapter.
The equal mass constraint also implies that the opening angles should be nearly the same, hence $|\theta_{ij}^{\text{open}} - \theta_{kl}^{\text{open}}| < 15^\circ$ is required. Fig. 4.9 shows the distribution of the difference in opening angles.

Figure 4.9: The difference in opening angles for multihadronic events satisfying aplanarity $> 0.01$ and $\bar{M} = 30.0 \pm 2.5$ GeV, and minimizing Eq. 4.9. Events are accepted if they satisfy $|\theta_{ij}^{\text{open}} - \theta_{kl}^{\text{open}}| < 15^\circ$. 
In contrast to $Z^0 \to q\bar{q}(g)$ decays, $Z^0 \to H^+H^-$ decays follow a $\sin^2\theta$ distribution, and therefore $\theta^p > 60^\circ$ is required. The production angle $\theta^p$ is shown in Fig. 4.10. Since the charged Higgs is a scalar, it decays isotropically in its rest frame. This is in contrast to gluon radiation from quarks, and therefore

$$\cos \theta^d_{ij} \cdot \cos \theta^d_{kl} < 0.13$$

(4.10)

is required. The quantity $\cos \theta^d_{ij} \cdot \cos \theta^d_{kl}$ is shown in Fig. 4.11.

The final-state topology varies widely as a function of Higgs mass, from being essentially two-jet-like with low sphericity and small average opening angle at low
masses to being nearly isotropic with high sphericity and large average opening angle at high masses. However, in the rest frame of each Higgs the topology is nearly two-jet-like with low sphericity throughout the accessible $H^\pm$ mass range since the Higgs masses are large relative to those of the daughter quarks. Consequently, the boosted sphericity product $S_{ij}^b \cdot S_{kl}^b$ [36] is required to satisfy

$$S_{ij}^b \cdot S_{kl}^b < 0.03.$$  (4.11)
Since the two charged Higgs particles have the same velocities, the boosted sphericities should not be vastly different, and therefore

$$0.067 < S_{ij}^b / S_{kl}^b < 15.0$$ \hspace{1cm} (4.12)

is also required. Fig. 4.12 shows the quantity $S_{ij}^b \cdot S_{kl}^b$ for multihadronic events with aplanarity $> 0.01$. Fig. 4.13 shows $S_{ij}^b \cdot S_{kl}^b$ for events which pass all the selection criteria except $S_{ij}^b \cdot S_{kl}^b < 0.03$.

Figure 4.12: The product of the boosted sphericities for multihadronic events satisfying aplanarity $> 0.01$ and $\bar{M} = 30.0 \pm 2.5$ GeV, and minimizing Eq. 4.9. Events are accepted if they satisfy $S_{ij}^b \cdot S_{kl}^b < 0.03$. 
Figure 4.13: The product of the boosted sphericities for events which pass all selection criteria except $S_{ij}^p \cdot S_{kl}^p < 0.03$. 
After all the selection criteria described above have been applied, six events remain in the data, one with an average reconstructed mass of \( \approx 25 \) GeV and the remaining five with average reconstructed masses \( > 40 \) GeV. The average reconstructed masses for the data, the \( Z^0 \rightarrow q\bar{q}(g) \) Monte Carlo, and a simulated 30 GeV charged Higgs are shown in Fig. 4.14. The disagreement be-

![Figure 4.14: The average reconstructed mass for events satisfying all charged Higgs selection criteria. The resolution \( \sigma_{M} \) using a Gaussian fit is 0.78 ± 0.06 GeV and the mean of the fit is 29.4 ± 0.1 GeV. The mean is low by 0.6 GeV due to the removal of tracks which do not satisfy the criteria for good tracks discussed at the beginning of this chapter.

between the number of data events (6) and the number of \( Z^0 \rightarrow q\bar{q}(g) \) Monte Carlo events (20 ± 2.2) satisfying all selection criteria is due to the fact that
the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo, for which 20 events represents only $\approx 0.02\%$ of the total number of Monte Carlo events generated, was not well-tuned for four-jet processes this far out on the tail of the various distributions. However, this disagreement is not a cause for great concern for the following reasons:

1. For the charged Higgs simulation, the fragmentation of the daughter quarks occurs at center-of-mass energies similar to those at PETRA, and the Monte Carlo has been well-tuned in this energy region [51]. Hence, the efficiency for the charged Higgs is not artificially increased like the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo.

2. In making the final limit contours, the estimated $Z^0 \rightarrow q\bar{q}(g)$ background is not used, i.e., no background subtraction is performed. Instead, only the observed data are used in making the limit, which is the more conservative technique [19].

The efficiencies for each of these selection criteria are given in Table 4.2 and the independent individual efficiencies in Table 4.3 for the data, the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo, and simulated charged Higgs signals at several values of $M_{H^{\pm}}$ for $H^+ H^- \rightarrow c\bar{s} \; \bar{c}s$. 
<table>
<thead>
<tr>
<th>Selection Criterion</th>
<th>$H^+ H^-$ Cumulative Efficiencies (%)</th>
<th>Z$^0 \rightarrow q\bar{q}(g)$ M.C.</th>
<th>$M_{H^\pm}$ = 22 GeV</th>
<th>$M_{H^\pm}$ = 30 GeV</th>
<th>$M_{H^\pm}$ = 35 GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>Eq. 4.9</td>
<td>49.8</td>
<td>47.0±1.4</td>
<td>93.0±6.8</td>
<td>72.4±3.5</td>
<td>45.2±3.0</td>
</tr>
<tr>
<td>Production Angle</td>
<td>25.1</td>
<td>24.2±1.0</td>
<td>68.7±5.8</td>
<td>54.5±3.1</td>
<td>31.2±2.5</td>
</tr>
<tr>
<td>Mass Difference</td>
<td>5.3</td>
<td>5.2±0.5</td>
<td>57.7±5.4</td>
<td>33.8±2.4</td>
<td>16.4±1.6</td>
</tr>
<tr>
<td>Opening Angles</td>
<td>1.5</td>
<td>2.1±0.3</td>
<td>46.8±4.8</td>
<td>24.1±2.0</td>
<td>11.8±1.5</td>
</tr>
<tr>
<td>Decay Angles</td>
<td>0.4</td>
<td>0.7±0.2</td>
<td>32.8±4.0</td>
<td>16.0±1.7</td>
<td>8.4±1.3</td>
</tr>
<tr>
<td>Boosted Sphericity</td>
<td>0</td>
<td>0.5±0.1</td>
<td>23.4±3.4</td>
<td>14.5±1.6</td>
<td>7.8±1.2</td>
</tr>
</tbody>
</table>

Table 4.2: The efficiencies of the remaining charged Higgs selection criteria after all previous criteria, including those of Table 4.1, have been applied, for the data, the background simulation, and simulated 22, 30, and 35 GeV charged Higgs. The data and Z$^0 \rightarrow q\bar{q}(g)$ M.C. efficiencies are calculated with the requirement $M = 30.0 \pm 2.5$ GeV, and each of the charged Higgs simulations with the requirement $M = M_{H^\pm} \pm 2.5$ GeV. The efficiencies are normalized to the number of events remaining after the application of the criteria of Table 4.1.
<table>
<thead>
<tr>
<th>Selection Criterion</th>
<th>H$^+$ H$^-$ Independent Efficiencies (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Data</td>
</tr>
<tr>
<td>Eq. 4.9</td>
<td>49.8</td>
</tr>
<tr>
<td>Production Angle</td>
<td>52.7</td>
</tr>
<tr>
<td>Mass Difference</td>
<td>12.0</td>
</tr>
<tr>
<td>Opening Angles</td>
<td>15.6</td>
</tr>
<tr>
<td>Decay Angles</td>
<td>25.3</td>
</tr>
<tr>
<td>Boosted Sphericity</td>
<td>81.0</td>
</tr>
</tbody>
</table>

Table 4.3: The efficiencies of the remaining charged Higgs selection criteria, taken independently, after the criteria of Table 4.1 have been applied. Otherwise, the same description as given in the caption of Table 4.2 applies.
4.2 Neutral SUSY Higgs Selection Criteria

Using the same definitions as in the charged Higgs analysis and the selection criteria listed in Table 4.1, the following additional criteria are applied to enhance the $h^0 A^0$ signal:

1. The production angle is required to satisfy $\theta_p > 60^\circ$;
2. The decay angles are required to satisfy $\max(\cos \theta^d_{ij}, \cos \theta^d_{kl}) < 0.8$;
3. The boosted sphericities are required to satisfy $\max(S^b_{ij}, S^b_{kl}) < 0.6$ and $0.067 < S^b_{ij}/S^b_{kl} < 15.0$.

The quantity $\max(\cos \theta^d_{ij}, \cos \theta^d_{kl})$ is shown in Fig. 4.15. The selection criteria used in the charged Higgs analysis which relied on the $M_{H^\pm} = M_{H^-}$ constraint—minimization of Eq. 4.9 and the differences in the reconstructed masses and opening angles—cannot be used for the $h^0 A^0$ analysis since the quantity $|M_{A^0} - M_{h^0}|$ can be large. The unequal masses also make the mass resolutions on $M_{A^0}$ and $M_{h^0}$ worse than that for $M_{H^\pm}$. The reconstructed mass for a 30 GeV $M_{A^0}$ is shown in Fig. 4.16. (The resolution on the reconstructed masses of the $h^0$ and $A^0$ are worse than that of the $H^\pm$ for several reasons. Chiefly, the resolution

---

1 Since the best of the three possible jet-jet pairings cannot be chosen by minimizing Eq. 4.9, in cases where two or all three of the pairings satisfy all selection criteria the pairing involving the higher energy jets is chosen. This pairing tends to correspond to reconstructed jet masses with larger values of $|M_{A^0} - M_{h^0}|$ than the other pairings, and as such populates the region near the edges of the limit contour. This is, therefore, a conservative approach in terms of making the limit.
Figure 4.15: The maximum of the cosines of the decay angles for multihadronic events satisfying aplanarity > 0.01 and $M_{h^0} = 30.0 \pm 5.0, M_{A^0} = 31.0 \pm 5.0$ GeV. Events are accepted if they satisfy $\max(\cos \theta^h_q, \cos \theta^A_q) < 0.8$.

on $M_{H^\pm}$ comes from average of two mass measurements, one from each equally-massed charged Higgs particle, whereas $M_{h^0}$ and $M_{A^0}$ must be measured independently. Also, the four $b$ quarks in the $h^0, A^0$ decay are more massive, have harder fragmentation characteristics, and decay more often into neutrinos than their counterparts in the $H^\pm$ decays, all of which make mass reconstruction for the neutral SUSY Higgs less accurate.) However, as indicated by Eq. 4.6, in the range of masses investigated in this thesis, both $h^0$ and $A^0$ decay into a pair of $b$ quarks, resulting in a significant number of leptons from semi-leptonic decays.
Figure 4.16: The reconstructed mass $M_{A^0}$ for multihadronic events satisfying aplanarity $> 0.01$. Since the best jet-jet pairing cannot, in general, be determined for the simulated $A^0, h^0$ signal, in the lower plot the Monte Carlo truth information is used and only the correct pairing is plotted. The resolution on $M_{A^0}$, using a Gaussian fit to the region between 26–36 GeV, is $2.9 \pm 0.2$ GeV.

of the $b$. The additional analysis leverage provided by these leptons permits a relaxation in the selection criteria (listed above) relative to the charged Higgs analysis. It also helps compensate for the lack of the equal-mass constraint.
Electrons from semi-leptonic b decays are identified with two independent measurements, the first from the pattern and magnitude of the energy deposition in the ECAL, and the second from the \( dE/dx \) measured by the TPC. These two measurements are complementary, with the first being more effective at high momenta \( (p > 5 \, \text{GeV}) \) and the second more effective at low momenta. Electron-positron pairs consistent with having come from a common vertex as the result of \( \pi^0 \) decay or photon conversion are removed from the sample by imposing the constraints that \( M_{e^+e^-} > 10 \, \text{MeV} \) and \( 5 \, \text{cm} < \rho < 40 \, \text{cm} \), where \( M_{e^+e^-} \) is the invariant mass of the pair and \( \rho \) is the radius of the common vertex. Muons from semi-leptonic b decays are identified by their penetration of, and minimal interaction with, a large number of the iron/Iarocci-tube layers of the HCAL. Identification of electrons and muons from semi-leptonic b decays are described in more detail in the following subsections. The lepton identification algorithms used here are the same as those used in Ref. 27.

4.2.1 Electron Identification

High momentum electrons measured by the ITC and TPC deposit most of their energy in only a few ECAL towers and in the earlier layers of the ECAL, and they lie on the relativistic plateau of the \( dE/dx \) curve as measured by the TPC. The main background to these electrons are charged pions, and in this momentum range they tend to deposit their energy over a larger region in the ECAL than electrons, and consequently are relatively easy to reject. Low momentum electrons \( (1 \, \text{GeV} < p < 5 \, \text{GeV}) \) are separated from low-momentum pions with
heavier reliance on the TPC $dE/dx$ measurement since the ECAL is considerably less effective in this momentum range.

The ECAL provides two measurements which are useful for identifying electrons. The first of these, $R_L$, compares the mean position of the shower profile depth (i.e., in an increasing radial direction) with that expected for an electron. The second, $R_T$, compares the quantity $E_4/p$ for a charged particle with that expected for an electron, where $E_4$ is the energy deposited in the four ECAL towers closest to the extrapolated position of the charged particle and $p$ is the momentum of the charged particle as measured by the tracking detectors. $R_T$ therefore takes advantage of both the fine granularity and the energy resolution of the ECAL.

The variable $R_L$ is constructed such that it is distributed as a standard normal Gaussian for electrons. The inverse of the mean shower position in depth is defined as

$$A = \frac{E_4}{\sum_{i=1}^{3} E_i S_i},$$

(4.13)

where $S_i$ is the mean shower position in depth for stack $i$. From test-beam data both $\langle A \rangle$ and $\sigma^2(A)$ are parametrized with energy and entry angle, and the quantity

$$R_L = \frac{A - \langle A \rangle}{\sigma(A)}$$

(4.14)

follows a standard normal distribution. A histogram of $R_L$ for electrons selected to come from photon conversions is shown in Fig. 4.17.
Figure 4.17: $R_L$ distribution for electrons from photon conversions. Photon conversions are selected by requiring $M_{e^+e^-} < 20$ MeV and $\rho > 4$ cm, where $M_{e^+e^-}$ is the invariant mass of the $e^+e^-$ pair and $\rho$ is the radius at which the photon converted. $R_L$ is plotted for one of the electrons if the other satisfies $-2.4 < R_L < 3.0$, $R_T > -3.0$ and $R_I > -2.5$ ($R_I$ is defined in Eq. 4.17).

Test beam data have shown that the variable

$$X = \frac{E_4}{p}$$  \hspace{1cm} (4.15)

also has a Gaussian distribution for electrons at a given energy. The mean, $\langle X \rangle$, is measured to be 0.83 from test-beam data, independent of angle and momentum provided $p > 2$ GeV. With the variance parametrized as $\sigma^2(X)$
using test beam data, for electrons the quantity

\[ R_T = \frac{X - \langle X \rangle}{\sigma(X)}, \tag{4.16} \]

also follows a standard normal distribution. A scatterplot of \( R_T \) for data satisfying \(-2.4 < R_L < 3.0 \) and \( p > 2 \text{ GeV} \) is shown in Fig. 4.18a.

![Scatterplot](image)

**Figure 4.18**: \( R_T \) and \( R_I \) distributions for the data. \( R_T \) is plotted for all events satisfying \(-2.4 < R_L < 3.0 \) as a function of momentum. Below \( p \approx 2 \text{ GeV} \), \( R_I \) no longer follows a standard normal distribution due to energy thresholds and entry angle effects. \( R_I \) is plotted for all events satisfying the \( R_L \) cut and \( R_T > -3.0 \).

The \( dE/dz \) measurement, \( I_m \), is defined as the 60% truncated mean of the individual TPC sense-wire measurements, and only tracks having at least 80 isolated sense-wire hits are considered. Hadronic data are used to fit the
\( \frac{dE}{dz} \) curves for pions, kaons, protons, muons and electrons, by parametrizing \( \sigma(I_m)/I_m \) with the number of wire samples. Although the \( \frac{dE}{dz} \) aids in electron identification at all but the highest momenta, it is especially useful at lower momenta where the \( R_L \) and \( R_T \) cuts are no longer so effective. For electrons, the mean \( \frac{dE}{dz} \) measurement \( \langle I_e \rangle \) lies a factor of 1.58 above the minimum, and the quantity

\[
R_I = \frac{I_m - \langle I_e \rangle}{\sigma(I_m)}
\]  

(4.17)

follows a standard normal distribution. A scatterplot of \( R_I \) for data satisfying \(-2.4 < R_L < 3.0 \) and \( R_T > -3.0 \) is shown in Fig. 4.18b. Electron candidates are required to satisfy \( R_I > -2.5 \).

The efficiency for electron identification for the ECAL alone is measured to be \( 80 \pm 2\% \), independent of \( p, p_t \). With the \( \frac{dE}{dz} \) measurement, the measured efficiencies are given in \( p, p_t \) bins in Table 1 of Ref. 27. They range between 37\% and 74\%. The table also gives the hadron misidentification probabilities for the ECAL alone and for the ECAL plus \( \frac{dE}{dz} \), which have values of 0.32–0.70\% and 0.02–0.25\%, respectively. The \( \frac{dE}{dz} \) measurement thus decreases the efficiency somewhat while greatly reducing the hadron misidentification probability.

### 4.2.2 Muon Identification

Muons typically make a minimum-ionizing energy deposit in the ECAL, and then penetrate many layers of the HCAL iron without undergoing significant multiple scattering. Usually, one Farocci tube per layer fires when a muon passes through. Each such interaction yields a "cluster," of which nearly the maximum
of 23, each lying very close to the extrapolation of the muon direction from the tracking detectors, are expected for muons. In contrast, pions which enter the HCAL tend to interact with the iron and leave large energy depositions over large areas of the HCAL, and therefore they do not leave the same linear pattern of fired tubes.

Muons are selected by extrapolating all charged tracks with $p > 1$ GeV into the HCAL and MCAL, where a cone three times larger than the r.m.s. multiple scattering width is formed around the track, and then requiring that

- more than 9 out of 23 planes fire in total,
- more than 4 of the last 10 planes fire,
- and at least 1 of the last 3 planes fire.

A plane is considered to have been fired by a muon provided its cluster consists of no more than 4 hits. Fig. 4.19 shows a typical pattern of fired planes left by a muon in the HCAL and muon chamber layers.

The hadronic contamination of the muon signal is small and due to three sources:

1. pion or kaon decay, with the resulting muon faking a muon from semi-leptonic $b$ decay,

2. "sail-through," when a hadron crosses the entire HCAL without interacting,
Figure 4.19: A typical pattern of fired planes left by a muon in the HCAL and muon chamber layers. The muon leaves a clear track in the ITC and TPC, a minimum ionizing deposit in the ECAL and HCAL, and a linear pattern of fired layers in the HCAL and muon chamber layers. The sizes of the energy deposits in the ECAL and HCAL are shown as histograms (with a bin width of one tower) and the scale for these deposits is shown in the upper right-hand corner of the diagram.

3. "punch-through," when an interacting hadron produces one or more secondaries which exit the HCAL through the multiple scattering cone.

The first two of these are well-predicted by the Monte Carlo, and contribute < 1% in total over the momentum range $p > 3$ GeV. To remove the punch-through background an additional cut is used in which a cone $\pm 25$ cm around the extrapolated track position is searched for excess hits and unassociated clusters.
Figure 4.20: Number of fired planes in the last ten of the HCAL for muon candidates which satisfy all other muon cuts listed on p. 82 for the data and $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo.

Fig. 4.20 shows the number of fired planes in the last 10 for muon candidates satisfying all other muon cuts.
The efficiencies for identifying muons are given in $p, p_t$ bins in Table 2 of Ref. 27. They range between 80% and 86% for $p > 3$ GeV. The table also gives the probability of contamination from $\pi^\pm$ or $K^\pm$ decay and the probability of hadron punch-through or sail through. These probabilities range between 0.21–1.10% and 0.29–0.56%, respectively.

4.2.3 Selection Criteria Based on Lepton Identification

Significant enhancement of the $h^0 A^0$ signal over the $Z^0 \to q\bar{q}(g)$ background is obtained by requiring, for all charged particles identified either as electrons with $p > 1$ GeV or muons with $p > 3$ GeV, that the following criteria be satisfied:

$$\sum_{\text{leptons}} p > 6.0 \text{ GeV}$$

and

$$\sum_{\text{leptons}} p_t > 0.75 \text{ GeV},$$

where $p_t$ for a given lepton is defined with respect to the closest jet. Figs. 4.21 and 4.22 show the $\sum p$ and $\sum p_t$ distributions, respectively, for the data, the $Z^0 \to q\bar{q}(g)$ Monte Carlo, and a simulated 30,31 GeV $h^0 A^0$ signal. The cumulative efficiencies of these selection criteria and the criteria listed on page 74 are given in Table 4.4 and the independent individual efficiencies in Table 4.5 for the data, the $Z^0 \to q\bar{q}(g)$ Monte Carlo, and simulated neutral SUSY Higgs signals at several values of $M_{h^0}$ and $M_{A^0}$. 
<table>
<thead>
<tr>
<th>Selection Criterion</th>
<th>h^0 A^0 Cumulative Efficiencies (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Data</td>
</tr>
<tr>
<td>Prod. Ang.</td>
<td>46.6</td>
</tr>
<tr>
<td>Σp</td>
<td>0.6</td>
</tr>
<tr>
<td>Σpt</td>
<td>0</td>
</tr>
<tr>
<td>Decay Angs.</td>
<td>0</td>
</tr>
<tr>
<td>Bstd. Sph.</td>
<td>0</td>
</tr>
</tbody>
</table>

Table 4.4: The efficiencies of the remaining neutral SUSY Higgs selection criteria after all previous criteria, including those of Table 4.1, have been applied, for the data, the background simulation, and simulated (20,35), (30,31), and (40,40) GeV neutral SUSY Higgs (h^0, A^0). The data and Z^0 \rightarrow q\bar{q}(g) M.C. efficiencies are calculated with the requirement (M_{h^0}, M_{A^0} = 30.0 \pm 5.0, 31.0 \pm 5.0) GeV, and each of the neutral SUSY Higgs simulations with the requirement that the reconstructed invariant masses are in the region (M_{h^0} \pm 5.0, M_{A^0} \pm 5.0) GeV. The efficiencies are normalized to the number of events remaining after the application of the criteria of Table 4.1.
Figure 4.21: The $\sum p$ distribution for multihadronic events satisfying aplanarity $> 0.01$, $M_{5\theta} = 30.0 \pm 5.0$ GeV, and $M_{A\theta} = 31.0 \pm 5.0$ GeV. Events are accepted if they satisfy $\sum p > 6.0$ GeV.
Figure 4.22: The $\sum p_t$ distribution for multihadronic events satisfying aplanarity $> 0.01$, $M_{h^0} = 30.0 \pm 5.0$ GeV, and $M_{A^0} = 31.0 \pm 5.0$ GeV. Events are accepted if they satisfy $\sum p_t > 0.75$ GeV.
<table>
<thead>
<tr>
<th>Selection Criterion</th>
<th>Data</th>
<th>( Z^0 \to q\bar{q}(g) )</th>
<th>( M_{h^0}, M_{A^0} = 20,35 \text{ GeV} )</th>
<th>( M_{h^0}, M_{A^0} = 30,31 \text{ GeV} )</th>
<th>( M_{h^0}, M_{A^0} = 40,40 \text{ GeV} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>Prod. Ang.</td>
<td>46.6</td>
<td>50.2±3.5</td>
<td>76.5±8.0</td>
<td>72.6±3.9</td>
<td>60.2±6.1</td>
</tr>
<tr>
<td>( \Sigma p )</td>
<td>3.4</td>
<td>4.5±1.0</td>
<td>37.0±5.6</td>
<td>29.1±2.4</td>
<td>27.3±4.1</td>
</tr>
<tr>
<td>( \Sigma p_t )</td>
<td>2.3</td>
<td>6.0±1.2</td>
<td>39.5±5.8</td>
<td>42.1±2.9</td>
<td>41.0±5.0</td>
</tr>
<tr>
<td>Decay Angs.</td>
<td>48.3</td>
<td>49.0±3.5</td>
<td>84.0±8.4</td>
<td>85.4±4.2</td>
<td>78.9±7.0</td>
</tr>
<tr>
<td>Bstd. Sph.</td>
<td>98.9</td>
<td>99.0±1.0</td>
<td>95.0±5.0</td>
<td>99.8±0.2</td>
<td>100.0±0.0</td>
</tr>
</tbody>
</table>

Table 4.5: The efficiencies of the remaining neutral SUSY Higgs selection criteria, taken independently, after the criteria of Table 4.1 have been applied. Otherwise, the same description as given in the caption of Table 4.4 applies.
After all the selection criteria described above have been applied, 22 data events and $19.8 \pm 2.5 \ Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo events remain. The reconstructed masses for the data, the $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo, and a simulated 30,31 GeV neutral SUSY Higgs pair are shown in Fig. 4.23. In contrast to the charged Higgs analysis, the agreement between the data and $Z^0 \rightarrow q\bar{q}(g)$ Monte Carlo is excellent because the Monte Carlo simulation of semi-leptonic decays is well understood and the cuts on identified leptons permit a relaxation in the other selection criteria.
Figure 4.23: The reconstructed masses $M_{h^0}$ and $M_{A^0}$ for events satisfying all neutral SUSY Higgs selection criteria for the 1989 data sample, $Z^0 \to q\bar{q}(g)$ Monte Carlo, and a simulated 30, 31 GeV $h^0 A^0$ pair. The solid lines demarcate the kinematically and theoretically allowed region for $(M_{h^0}, M_{A^0})$. The $Z^0 \to q\bar{q}(g)$ and neutral SUSY Higgs signals are normalized to the luminosity of the data.
Chapter 5

Experimental Results

The analyses presented in Chap. 4 show that there is no clear evidence in favor of the existence of pair-produced charged Higgs or neutral SUSY Higgs in the channels

\[
Z^0 \rightarrow \begin{cases} 
    H^+ H^- \rightarrow cs \bar{c}s, \\
    H^+ H^- \rightarrow cb \bar{c}b \text{ or} \\
    h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}.
\end{cases}
\]  (5.1)

This is most evident from Figs. 4.14 and 4.23, where the correspondence between data and \(Z^0 \rightarrow q\bar{q}(g)\) Monte Carlo is much closer than that between data and any Higgs signal. Hence, a 95% confidence level (CL) limit is set for each of the processes above. The procedure by which this is done is described in Sect. 5.2. Before forming the limits, however, several systematic checks are performed to insure that the limit is not unstable to changes in the parameters used in the simulation, changes in the selection criteria, or statistical fluctuations leading to an incorrect estimation of the signal efficiency.
5.1 Systematics

It is crucial that the Higgs signal processes are simulated with sufficient accuracy. Otherwise, the selection criteria designed to enhance these processes might instead remove genuine signal events from the data or incorrectly classify $Z^0 \to q\overline{q}(g)$ events as Higgs events. The most important variables in this regard are the LUND fragmentation parameters (described in detail in Sect. 3.2.1). Although the fragmentation parameters used here for the Higgs signal simulation have been well-tuned to PEP and PETRA data [51] (where the center-of-mass energies were comparable to the scalar rest masses relevant here), it is nevertheless essential to check what effect a reasonable change in the fragmentation parameters has on the signal efficiency.

To do this, variations in the LUND parameters are made, and several sets of 30, 30 GeV $h^0 \ A^0$ events are generated. These sets of events, one per variation, are then passed through the full simulation and reconstruction codes (GALEPH and JULIA, respectively), and finally all the selection criteria discussed in Chapter 4 are applied. The efficiencies for signal detection are then calculated for each set of events. The LUND parameters are varied over large ranges as shown in Table 5.1. A similar procedure is performed for the charged Higgs as well.

In no instance does the signal detection efficiency change by more than 10% from its original value. The possible effect of this change is taken into account in a conservative manner by reducing the signal efficiency by the corresponding amount when making the limit.
<table>
<thead>
<tr>
<th></th>
<th>Original Value</th>
<th>Range of Variation</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\sigma_{x,y}$</td>
<td>0.325</td>
<td>0.300–0.400</td>
</tr>
<tr>
<td>$a$</td>
<td>0.45</td>
<td>0.30–1.00</td>
</tr>
<tr>
<td>$b$</td>
<td>0.90</td>
<td>0.65–1.10</td>
</tr>
<tr>
<td>$\epsilon_c$</td>
<td>0.025</td>
<td>0.015–0.075</td>
</tr>
<tr>
<td>$\epsilon_b$</td>
<td>0.015</td>
<td>0.008–0.030</td>
</tr>
</tbody>
</table>

Table 5.1: The range of variation of the LUND parameters $\sigma_{x,y}$, $a$, $b$, $\epsilon_c$, and $\epsilon_b$ (see Sect. 3.2.1).

The stability of the final limit to variations in the values of the selection criteria is checked by varying the criteria within reasonable levels and then evaluating the resulting change in the Higgs signal and data efficiencies. A variation in the selection criteria of $\pm 10\%$ produces virtually no change in either of these. Nevertheless, an estimated $5\%$ error is assumed and taken into account as described above.

A small $2\%$ systematic error in the luminosity measurement [37] is also taken into account. Finally, since the signal processes are generated in finite numbers, a statistical fluctuation could result in a systematic error in the calculation of the signal efficiency. The error is estimated by artificially varying the efficiencies measured for various Higgs masses within their Gaussian errors, and then refitting the efficiency curve (see Sect. 5.2) to these artificially varied points. The
measured variation in the efficiency derived from this curve at several arbitrarily chosen Higgs masses indicates that this error is less than 10%.

Adding in quadrature all the sources of error listed above gives a total systematic error of 15%. This is the amount by which the expected Higgs signals are reduced before the final limits are made.

5.2 Calculating the Limits

The lack of evidence for any of the decay modes of eq. 5.1 is quantified probabilistically. This quantification is done by assuming that the signal is present, and then, from the number of events in the data which satisfy all selection criteria, estimating the level of its assumed presence in a conservative manner. Since the sample of events to which the criteria are applied is large, and the number satisfying them is small, the latter are taken to follow a Poisson distribution. The estimation of the level of the presence of the signal in the data \( N \) is taken conservatively as that value of \( N \) such that it would be exactly 95% probable that, on average, repeated experiments would see more than the number observed in this experiment \( n_0 \). In other words, the \( n_0 \) are assumed to come from a Poisson distribution of mean \( N \) such that in 95% of repeated experiments more than \( n_0 \) events would have been observed. Since this number \( N \) is the assumed signal level in the data, only when the expected signal level from simulations of the Higgs processes \( \mu_* \) exceeds \( N \) at some \( M_{H^\pm} \) is it stated that the charged Higgs of mass \( M_{H^\pm} \) is excluded at 95% CL. The estimations of \( \mu_* \)
<table>
<thead>
<tr>
<th>$n_0$</th>
<th>0</th>
<th>1</th>
<th>2</th>
<th>3</th>
<th>4</th>
<th>10</th>
</tr>
</thead>
<tbody>
<tr>
<td>$N$</td>
<td>3.00</td>
<td>4.74</td>
<td>6.30</td>
<td>7.75</td>
<td>9.15</td>
<td>16.96</td>
</tr>
</tbody>
</table>

Table 5.2: Several Poisson upper limits $N$ at CL = 95% for $n_0$ observed events.

for the charged Higgs and the neutral SUSY Higgs are discussed in Sects. 5.2.1 and 5.2.2, respectively.

The Poisson distribution with mean $\mu$ is:

$$f(n; \mu) = \frac{e^{-\mu} \mu^n}{n!}, \quad n = 0, 1, 2, \cdots$$

(5.2)

where $n$ is the number of observed events. The confidence level is:

$$\text{CL} = \sum_{n=n_0+1}^{\infty} f(n; N)$$

$$= 1 - \sum_{n=0}^{n_0} f(n; N)$$

(5.3)

where $n$ is the number of events seen in a repeated experiment, $n_0$ and $N$ are as defined above, and CL = 0.95. Some sample values of $n_0$ and $N$ at 95% CL are given in Table 5.2 [19].

5.2.1 Limits on Charged Higgs Production

In the mass range of the charged Higgs explored in this thesis, the decay modes

$$Z^0 \rightarrow \begin{cases} 
H^+ \ H^- \rightarrow \nu \bar{\nu} \ \bar{\tau} \tau, \\
H^+ \ H^- \rightarrow \nu \bar{\nu} \ \bar{c} \bar{s} \quad \text{and} \quad H^+ \ H^- \rightarrow \nu \bar{\nu} \ \bar{c} \bar{b}
\end{cases}$$

(5.4)

can occur in addition to the four-jet modes

$$Z^0 \rightarrow \begin{cases} 
H^+ \ H^- \rightarrow c \bar{s} \ \bar{c} \bar{s} \quad \text{and} \\
H^+ \ H^- \rightarrow c \bar{b} \ \bar{c} \bar{b}.
\end{cases}$$

(5.5)
Since the relative rates of these processes are not predicted by the theory, the limit on the pair production of charged Higgs is expressed in terms of the branching ratios \( \text{BR}[H^\pm \to cs] \) and \( \text{BR}[H^\pm \to cb] \). Defining \( B_h \equiv \text{BR}[H^\pm \to cs] \) and \( B_l \equiv \text{BR}[H^\pm \to \nu \tau] \), the probability that both charged Higgs decay hadronically is \( B_h^2 \), the probability that one Higgs decays hadronically and the other leptonically is \( 2B_hB_l \), and the probability that both decay leptonically is \( B_l^2 \).

If \( \mu_n \) charged Higgs events are expected at \( B_h = 100\% \), then \( B_h^2 \cdot \mu_n \) events in which both Higgs decay hadronically should be seen, so the 95\% CL limit on the branching ratio is obtained by solving the equation

\[
B_h^2 \cdot \mu_n = N
\]  

(5.6)

for \( B_h \), where \( N \) is as defined above. The same argument applies to the mode \( H^+ H^- \to \bar{c}b \; c\bar{b} \), which is treated independently of the mode \( H^+ H^- \to c\bar{s} \; \bar{c}s \).

The searches in the decay modes shown in Eq. 5.4, along with the four-jet modes discussed in this thesis, are presented in detail in Ref. 50. The 95\% CL limits derived by these searches are discussed below.

The number of expected charged Higgs events before the application of any selection criteria is calculated from the measured luminosity and the predicted cross section given by Eq. 1.19. The detection efficiencies are calculated for simulated charged Higgs events generated at several different values of \( M_{H^\pm} \) and these efficiencies are then fit to a polynomial [52], as shown in Figs. 5.1 and 5.2 for the two different four-jet decay modes under study. At intervals in \( M_{H^\pm} \) of 0.1 GeV over the mass range 10-40 GeV, the efficiency is estimated from the
polynomial fit and the number of expected charged Higgs events after the application of all selection criteria is calculated at each point. To make the limit on $\text{BR}[H^\pm \rightarrow cs(c\bar{b})]$, only events from the region $\pm 2.5$ GeV ($\approx \pm \sigma_M$, the resolution on the reconstructed mass before the application of the selection criteria of Sect. 4.1) around each point are used. For example, if only 1 event from the data satisfies all the selection criteria, and it has an average reconstructed mass $\bar{M}_{\text{data}}$, Eq. 5.6 becomes:

$$B_h^2 \cdot \mu_s = \begin{cases} 4.74 & \text{for } \bar{M} = \bar{M}_{\text{data}} \pm 2.5 \text{ GeV} \\ 3.00 & \text{otherwise.} \end{cases} \tag{5.7}$$

In the analysis presented here, $\bar{M}_{\text{data}} \approx 25$ GeV.

The 95% CL limit on $\text{BR}[H^\pm \rightarrow cs]$ is shown in Fig. 5.3, and the limit on $\text{BR}[H^\pm \rightarrow c\bar{b}]$ in Fig. 5.4. In both plots the expected Higgs signals are reduced by 15% to account for the systematic errors discussed in Sect. 5.1. The limit on $\text{BR}[H^\pm \rightarrow c\bar{b}]$ is slightly weaker than that for $\text{BR}[H^\pm \rightarrow cs]$ due primarily to the lower efficiency of the boosted sphericity product selection criterion. This is because the boosted sphericity is larger in the mode $H^+ H^- \rightarrow c\bar{b} c\bar{b}$ since the $b$ quark has a larger mass and a harder fragmentation than the $s$ quark. In the low $M_{H^\pm}$ region, the charged Higgs events become very two-jet-like, making it difficult to distinguish them from the copious two-jet $Z^0 \rightarrow q\bar{q}(g)$ background and preventing the limit from going lower. Also, in this region the boosted sphericity product selection criterion becomes inefficient because the lighter Higgs do not give a large boost to their decay products. Importantly, however, the lower end of the limit passes well below the previous limit of 19 GeV [23]. In the
high $M_{H^\pm}$ region, the decrease in production rate and the difficulty in accurately reconstructing four jets are the two main factors which prevent the limit from going higher. For the same reasons, the efficiency curve has the shape shown in Fig. 5.1. Results similar to those presented here have been obtained by other LEP collaborations [53].
Figure 5.1: The fit to the detection efficiencies for the charged Higgs decaying in the mode $H^+ H^- \rightarrow cs \, \bar{c}s$. A $\chi^2$ fit to a fourth order Chebyshev polynomial is used. The chi squared per degree of freedom for the fit is 0.2. At $M_{H^\pm} = 13.0$ GeV the detection efficiency is arbitrarily set equal to 0.1% in order to keep the fit well-behaved. The shape of the curve is explained in the text.
Figure 5.2: The fit to the detection efficiencies for the charged Higgs decaying in the mode $H^+ H^- \rightarrow c \bar{b} \bar{c} b$. A $\chi^2$ fit to a fourth order Chebyshev polynomial is used. The chi squared per degree of freedom for the fit is 1.43. At $M_{H^\pm} = 13.0$ GeV the detection efficiency is arbitrarily set equal to 0.1% in order to keep the fit well-behaved. The shape of the curve is explained in the text.
Figure 5.3: The 95% CL limit on $\text{BR}[H^\pm \to cs]$ as a function of $M_{H^\pm}$ after a 15% reduction in the expected signal to account for systematic uncertainties. The notch centered at $\approx 25$ GeV is due to a data event which satisfied all the selection criteria and which had an average reconstructed mass of $\approx 25$ GeV. The excluded region extends from 14.5–35.2 GeV at $\text{BR}[H^\pm \to cs] = 100\%$. Previous work excludes the charged Higgs below 19 GeV [23].
Figure 5.4: The 95% CL limit on $\text{BR}[H^\pm \to cb]$ as a function of $M_{H^\pm}$ after a 15% reduction in the expected signal to account for systematic uncertainties. The notch centered at $\approx 25$ GeV is caused by the data event which satisfied all the selection criteria and had an average reconstructed mass of $\approx 25$ GeV. The excluded region extends from 16.4–35.4 GeV at $\text{BR}[H^\pm \to cb] = 100\%$. Previous work excludes the charged Higgs below 19 GeV [23].
Analyses designed to detect the charged Higgs decay modes

\[ Z^0 \rightarrow \begin{cases} 
H^+ H^- \rightarrow \nu \overline{\nu} \tau, \\
H^+ H^- \rightarrow \nu \overline{\nu} \bar{c}s \quad \text{and} \quad H^+ H^- \rightarrow \nu \overline{\nu} \bar{c}b,
\end{cases} \quad (5.8) \]

have also shown no clear evidence of the charged Higgs, and so 95\% CL limits on \( BR[H^\pm \rightarrow \nu \tau] = 1 - BR[H^\pm \rightarrow cs(c\bar{b})] \) have been derived. These limits, along with the limits derived for the four jet decay mode, are shown in Figs. 5.5 and 5.6 for the \( cs \) and \( cb \) hadronic decay modes, respectively. Also shown in these figures is the combined 95\% CL limit for all three processes, obtained by solving the equation

\[ B_l^2 \cdot \mu_{ll} + 2(1 - B_l)B_l \cdot \mu_{lh} + (1 - B_l)^2 \cdot \mu_{hh} = N_{ll} + N_{lh} + N_{hh} \quad (5.9) \]

for \( B_l \), where \( B_l = BR[H^\pm \rightarrow \nu \tau] \), \( \mu \) is the expected signal \( N \) is the number of observed events, and the subscripts \( ll \), \( lh \) and \( hh \) stand for the decay modes

\[ H^+ H^- \rightarrow \nu \overline{\nu} \tau, \]

\[ H^+ H^- \rightarrow \nu \overline{\nu} \bar{c}s \quad \text{and} \]

\[ H^+ H^- \rightarrow c\bar{s} \bar{c}s, \quad (5.10) \]

respectively. The same procedure was applied for the decay mode \( H^+ H^- \rightarrow cb \bar{c}b \). In regions not contained by any of the three individual 95\% CL contours on \( BR[H^\pm \rightarrow \nu \tau] \), four possible combined limits on \( BR[H^\pm \rightarrow \nu \tau] \) are considered: one with \( \mu_{ll} \) set to zero, one with \( \mu_{lh} \) set to zero, one with \( \mu_{hh} \) set to zero, and one with all three \( \mu \)'s having their measured values. The combination yielding the best limit on \( BR[H^\pm \rightarrow \nu \tau] \) is taken as the combined limit at a given \( M_{H^\pm} \).
Figure 5.5: The 95% CL limit on $\text{BR}(H^\pm \rightarrow \nu \tau)$ and $\text{BR}(H^\pm \rightarrow cs)$ as a function of $M_{H^\pm}$ after a 15% reduction in the expected signal to account for systematic uncertainties. The region between the dashed lines represents the combined limit. Previous work excludes the charged Higgs below 19 GeV [23].
Figure 5.6: The 95% CL limit on BR($H^\pm \rightarrow \nu\tau$) and BR($H^\pm \rightarrow cb$) as a function of $M_{H^\pm}$ after a 15% reduction in the expected signal to account for systematic uncertainties. The region between the dashed lines represents the combined limit. Previous work excludes the charged Higgs below 19 GeV [23].
5.2.2 Limits on Neutral SUSY Higgs Production

In the mass range of the neutral SUSY Higgs explored in this thesis, (both $M_{h^0}$ and $M_{A^0} > 15$ GeV) the four-jet decay mode

$$Z^0 \rightarrow h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}$$

(5.11)

dominates over all other possible modes. There is always a contribution from the decay mode

$$Z^0 \rightarrow h^0 A^0 \rightarrow \tau\bar{\tau} \ b\bar{b},$$

(5.12)

but since this is only at the 5% level the limit on the four-jet mode is made assuming $\text{BR}[h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}] = 100\%$. The results of the search in the decay mode shown above in Eq. 5.12, along with other decay modes and the four-jet mode discussed in this thesis, are presented in Ref. 55.

The number of expected neutral SUSY Higgs events before the application of any selection criteria is calculated from the measured luminosity and the predicted cross section given by Eq. 1.20. The detection efficiencies are calculated for simulated neutral SUSY Higgs events generated at 19 different points in the $M_{h^0}$, $M_{A^0}$ plane, and these efficiencies are then fit to a two-dimensional surface using a weighted least-squares bicubic spline fit [54]. These efficiencies, and the corresponding fit, are shown in Fig. 5.7 for various $M_{h^0}$ as a function of $M_{A^0}$. At intervals in $M_{h^0}$ and $M_{A^0}$ of 0.5 GeV over the mass ranges $15.0$ GeV $< M_{h^0} < 48.0$ GeV and $15.0$ GeV $< M_{A^0} < 48.0$ GeV, the efficiency is estimated from the surface fit and the number of expected charged
Higgs events after the application of all selection criteria is calculated at each point. To make the limit on $M_{h^0}$ and $M_{A^0}$, only events from the region $\pm 5.0$ GeV ($\approx \pm \sigma_M$, the resolution on either reconstructed mass before the application of the selection criteria of Sect. 4.2) around each point are used.

The 95% CL limit on $(M_{h^0}, M_{A^0})$ is shown in Fig. 5.8. In this plot the excluded region lies between the curve and the line described by $M_{h^0} = M_{A^0}$. For reference, the line corresponding to the kinematic limit at LEP-I energies is also shown. The limit derived from the minimal Standard Model Higgs search [17] excludes a region delineated by the dashed line in the plot. Results from the other channels are presented in Ref. 55. As with the limit contours for the charged Higgs, the expected neutral SUSY Higgs signals are reduced by 15% to account for the systematic errors discussed in Sect. 5.1.
Figure 5.7: The fit to the detection efficiencies for the neutral SUSY Higgs decaying in the mode $h^0 A^0 \rightarrow b\bar{b} \ b\bar{b}$ for several values of $M_{h^0}$. The efficiencies are fit to a two-dimensional surface using a weighted least-squares bicubic spline. At $M_{A^0} = 55.0$ GeV the detection efficiency is arbitrarily set equal to 0% in order to keep the fit well-behaved.
Figure 5.8: The 95% CL limit on the production of $M_{h^0}, M_{A^0}$ after a 15% reduction in the expected signal to account for systematic uncertainties. The region between the solid curve and the line described by $M_{h^0} = M_{A^0}$ is excluded by the analysis of Sect. 4.2. This region extends from 15.0–40.0 GeV along the line $M_{h^0} = M_{A^0}$, and as high as $M_{A^0} = 44.0$ GeV at $M_{h^0} = 24.0$ GeV. The region to the left and above the dashed line is excluded by the minimal Standard Model Higgs search. The region below $M_{h^0} = 15.0$ GeV and $M_{A^0} = 20.0$ GeV is excluded by other searches described in Ref. 55.
5.3 Conclusion

Searches for hadronically-decaying, pair-produced charged Higgs bosons and pair-produced neutral SUSY Higgs bosons at the ALEPH detector at LEP have not shown any evidence for the existence of their signatures. Limits on the masses of these particles within the frameworks of the two-doublet minimal Standard Model and the minimal supersymmetric model, respectively, have therefore been made. The search for the charged Higgs relied on the constraint provided by the equal masses of the $H^+$ and $H^-$ to obtain a mass resolution of $\approx 2.0$ GeV throughout the mass range of interest, and on the incremental contributions of many cuts to enhance the signal-to-noise ratio to levels at which a charged Higgs would have been visible. Since the charged Higgs can decay into modes other than the four-jet mode investigated here, with branching ratios which are free parameters of the theory, the limit is expressed in terms of this branching ratio, $\text{BR}[H^\pm \to cs(cb)]$. The 95% CL limit on the charged Higgs extends from 14.5–35.2 GeV at $\text{BR}[H^\pm \to cs] = 100\%$, and from 16.4–35.4 GeV at $\text{BR}[H^\pm \to cb] = 100\%$.

The search for neutral SUSY Higgs bosons was made in a mass region where the dominant decay mode yields four $b$ quarks in the final state. Thus, the signature is also four jets like with the charged Higgs, but unlike the charged Higgs the neutral SUSY Higgs bosons can be produced with unequal masses, removing an important constraint and significantly degrading the mass resolution. However, with four $b$ quarks in the final state a significant number of high $p_T$
leptons from semi-leptonic decays of the $b$ would be expected, and this signature provided a crucial improvement in the signal-to-noise ratio. Since the masses of the neutral SUSY Higgs bosons are not necessarily equal, the limit is expressed as an excluded region in the $M_{h^0}, M_{A^0}$ mass plane. The 95% CL limit on these particles extends from 15.0–40.0 GeV along the line $M_{h^0} = M_{A^0}$, and as high as $M_{A^0} = 44.0$ GeV at $M_{h^0} = 24.0$ GeV.
Appendix A

The TPC Gating System

A.1 Introduction

The operation of a TPC in a continuously sensitive mode can pose a severe problem: the buildup of space charge in the drift region. This space charge is comprised of positive ions which are mainly produced during the proportional amplification at the sense wires and which subsequently enter the drift region. Once in the drift region, the space charge can alter the local electric field, causing unwanted track distortions. For this reason, it was decided at an early stage in the design of the ALEPH TPC to implement an extra grid of wires which would act to prevent positive ions from entering the drift region.

A.2 Principle of Gating

The extra grid of wires, or “gating grid” [56] is installed between the shielding grid and the drift region, as shown in fig A.1.

The gating grid is either in the “open” state or the “closed” state. In the open state (Fig. A.1a) the potential $V_g (\sim -67V)$ is placed on the gate wires
so that the gate is transparent to the passage of drifting charged particles. In the closed state (Fig. A.1b), the potentials $V_g \pm \Delta V_g$ are placed on alternate wires of the grid so that the resulting dipole fields render the gate opaque to the passage of charged particles. A $\Delta V_g$ of $\leq 40V$ suffices to block the passage of the positive ions [57] while, because of the magnetic field, a larger $\Delta V_g$ of $\geq 200V$ is required block electrons, as shown in Fig. A.2.

Switching between the closed and open state of the gate is done synchronously with the bunch crossing (BX). This method of gating is illustrated in Fig. A.3. About 3 $\mu$s before the BX the gate is opened to allow electrons to drift into the amplification region. Normally the gate is then closed when the first level trigger is negative (about 3 $\mu$s after the BX). Only when the first level trigger is positive is the gate held open for the maximum 45 $\mu$s drift time of the electrons in the TPC.

It has been shown [57] that even though this synchronous mode of operation leaves the gate open roughly 25% of the time positive ions are still unable to penetrate the gating grid. The positive ion transparency of a synchronously operated gating grid as a function of the beam-crossing frequency is shown in Fig. A.4. The gate remains effective because the positive ions do not drift very far during one gating cycle and can still be neutralized at the gating grid, after following a path such as the one depicted in Fig. A.5. This mode of gating has been named "synchronous ion trapping."

As shown by Fig. A.4, $\Delta V_g \approx 50V$ would more than sufficient to keep the gate opaque to positive ions even in synchronous operation. However, at
such a low $\Delta V_g$ the gate would still be rather transparent to incoming electrons. By using a higher voltage ($\Delta V_g = 100V$) the closed gate becomes about 50% more opaque to electrons. This improved ability to keep out electrons from background sources (primarily synchrotron radiation and beam-gas interactions) lengthens the chambers’ lifetime. A further reduction in electron transparency can be obtained by using an even higher $\Delta V_g$ at the expense of higher pickup at the sense wires when the gate switches. The option of keeping the gate either closed or open for indefinite periods of time is also available.

A.3 Design Considerations

The gating circuit has been designed not only to place large voltage swings on the gating grid according to the synchronous ion trapping scheme, but also to minimize the pickup (primarily at the sense wires) resulting from these swings. Pickup is particularly worrisome when the gate opens since any significant pickup which still exists after bunch crossing would corrupt incoming data. Pickup could also produce too much extra data on the wires, which could overload the processors unless the threshold were raised. If the pickup at gate opening or closing were so large that it saturated the shapers, an unpredictable DC offset in the shapers could result (see Sect. A.4). Large pickup could also cause problems for the trigger pad electronics. To minimize the pickup, the gating circuit allows one to tune the (+) and (−) gating waveforms, that is, to match the shapes of the pulses going to the (+) and (−) gating wires in order to make them as identical as possible. For the opening and, to a lesser extent, the closing waveforms,
the circuit can be tuned by adjusting several points within the circuit. A block diagram of the gating circuit is shown in Fig. A.6. There is one circuit per sector in order to achieve ground isolation on a sector-by-sector basis.

The power supplies needed for the operation of these circuits are physically located in close proximity to the circuit boards themselves. There is one power supply per circuit board (and per sector), also for ground isolation purposes.

### A.4 Performance

At $\Delta V_g = 150\text{V}$, the wires have negligible pickup 3 $\mu$s after gate opening and the opacity to electrons is enhanced significantly over $\Delta V_g = 100\text{V}$. However, such a high $\Delta V_g$ makes sectors more susceptible to excessive pickup at gate closing, which is difficult to remove by tuning the circuit since the circuit was designed mainly to match the opening waveforms. During the 1989 data run, when $\Delta V_g = 150\text{V}$, certain sectors experienced pickup induced by the gate closing which saturated the electronics and resulted in a pedestal shift, adversely affecting the $\text{d}E/\text{d}x$ measurement. This problem has been remedied by reducing the $\Delta V_g$ to 100$\text{V}$. Recent measurements on the 1990 data have shown that the pickup induced by the gating no longer affects the $\text{d}E/\text{d}x$ measurements, and is below normal TPD thresholds. In summary, the gating has no detrimental effect on the data taken in the TPC, prevents the buildup of space charge and, by partially sealing off the proportional region for about 40% of each cycle, extends the lifetime of the TPC chambers.
Figure A.1: The TPC gating grid, located 14 mm from the pad plane, in the open (a) and closed (b) states. In the open state, all the wires of the gating grid are held at the ambient potential of the drift field; in the closed state, additional voltages are placed on alternating wires to set up the dipole fields shown above. Just above the pad plane is the sense/field wire grid, where the proportional amplification occurs. Between the sense/field wire and gating grids is the cathode (or ground) grid, which separates the high field region in the vicinity of the sense wires from the drift field region.
Figure A.2: The transparency of the TPC gating grid for electrons in the presence of a magnetic field as a function of $\Delta V_g$. The transparency for positive ions is nearly independent of magnetic field and hence coincides with the one for electrons at $B = 0$. 
Figure A.3: The synchronous mode of gating grid operation.
Figure A.4: The positive ion transparency of a synchronously operated gating grid as a function of the beam crossing frequency. In this plot, the grid is left open for 6μs, independent of frequency.
Figure A.5: An illustration of how a positive ion is neutralized by the gating grid.
Figure A.6: A block diagram of the gating circuit.
Appendix B

Compensation for Edge Effects in TPC Sectors

B.1 Introduction

At distances up to 1 cm from the edge of the TPC sectors significant position and gain distortions have been observed. Measurements made using a small chamber [58] showed conclusively that three strips placed at the chamber edge (see Fig. 2.6) could effectively correct for position and gain distortions at distances of up to 1 cm from the sector edge. The measurement of gain distortion, defined as the variation of the gain as a function of the distance from the sector edge relative to that measured at distances far from the edge, was performed with K and W sectors (inner and outer sectors, respectively; see Fig. 2.5) using the sector test stand set-up, which was originally constructed for making gain maps of the sectors, at MPI (Max Planck Institut) in Munich. A schematic of this set-up is shown in Fig. B.1. It was not possible to do position distortion measurements since the stand used an \(^{55}\text{Fe}\) source which could not be sufficiently
Figure B.1: A schematic of the test stand set-up in Munich, showing the K-type sector, the movable $^{55}$Fe source, a pad row, and a few sense wires.

collimated for such a measurement. However, edge position distortions are expected to be small in ALEPH’s 1.5 T magnetic field since $\omega \tau \simeq 7$ in this field. For the remainder of this appendix only gain distortion measurements will be discussed.

The appropriate potential strip voltage was found to be a roughly linear function of sense and field wire voltage, and it also depended on the angle $\theta$ the sense wires made with the edge. Over a large portion of the edge region, $\theta$ differs for K and W sectors, but it is the same for W and M sectors. The $\theta$
dependence suggests that the TPC ought to have different strip voltages for K sectors relative to M and W sectors.

B.2 Apparatus

The MPI test stand (see Fig. B.1) consisted mainly of a sector mount, a gas system providing Ar:CH₄ (90 : 10) at 1 atmosphere as in the ALEPH TPC, voltage supplies and a movable $^{55}$Fe strip source. All wire channels were equipped for readout via standard TPC preamps and shapers. Commercial ADC’s were used to digitize the signal.

The voltages on all three potential strips, the gating grid, and the sense and field wires were independently adjustable. Note that $V_{\text{upper}}$, the potential on the strip between the gating and ground grids, was held at $(V_{\text{gate}} + V_{\text{ground}})/2$, in accordance with the results of small chamber tests which indicated that this potential would minimize position distortions [58]. The thin aluminized mylar sheet, positioned about 2 cm from the gating grid plane, was held at ground to eliminate charge accumulation in the region between the gating grid and the outer boundary of the test chamber gas volume. Experience with gain maps has shown that charge accumulation in this region has adverse effects on the measured gain. The source was collimated to ensure no more than a 2 mm spread in the illumination at the sense wires.

Gamma rays emitted by the $^{55}$Fe source could convert in the gas anywhere in the region between the aluminized mylar sheet and the pad plane but, by holding the gating grid at a uniform negative potential, only those gammas converting
between the gating grid and the pad plane were detected. Conversions outside of this region were intentionally excluded because the electric field there could not be well-defined. The accepted conversions between the gating grid and pad plane, however, were different in one very important characteristic from those accepted by the TPC in normal operation. In the TPC, ionization nearly always arrives at the sense wires from "above", i.e., after passing through the gating and ground grids. For this measurement, however, roughly 25% of the ionization arrived at the sense wire from "below", i.e., was deposited between the pad and sense-field wire planes. This raised the concern that the measurement would be adversely affected if there were different gains for ionization deposited above and below the sense wires.

To determine the significance of this effect, a method was devised to make the gain distortion measurement using only ionization which came from above the sense wires. Then, of course, only the gain experienced by clusters impinging on the sense wires from above would be registered. The essential idea was to use the gating grid to set up two different drift fields: the first, already mentioned above, was obtained by placing a uniform negative voltage on the gating grid; the second by using a uniform positive voltage. In the first configuration, ionization from the region between the gating grid and the pad plane could impinge on the sense wires. In the second configuration, ionization from between the ground grid and the pad plane could reach the sense wires, and any ionization deposited between the gating grid and ground grid drifted back to the gating grid where it was neutralized. A bin-by-bin subtraction of the $^{55}$Fe spectra obtained with the
positive voltage from that obtained with the negative voltage gave a spectrum which consisted statistically of ionization deposited only in the region between the gating grid and ground grid, i.e., from above the sense wires.

Fortunately, the gain distortions measured using this subtraction technique were virtually identical to those measured using just a uniform negative or positive gating grid voltage. Thus the duration of each data taking run was shortened, and, more importantly, in future operation there is no need to be concerned with differences in gain from ionization deposits impinging the sense wires from above as opposed to below.

B.3 Procedure

The $^{55}$Fe source was moved to the center of the sector as shown in Fig. B.1. The well-known $^{55}$Fe spectrum was collected for each wire by histogramming the peak pulse height produced by a wire pulse. A sample spectrum is shown in Fig. B.2. The mean of a gaussian curve fitted to the main peak of this spectrum $\hat{x}_{\text{center}}$, was found using a least-squares fit. This mean corresponds directly to the gain. Next, the source was moved to the edge where the mean of the spectrum peak, $\hat{x}_{\text{edge}}$, was similarly obtained.

For each wire, the relative gain was then calculated:

$$\frac{\hat{x}_{\text{edge}}}{\hat{x}_{\text{center}}}$$  \hspace{1cm} (B.1)

Normalizing each wire to itself in this manner effectively eliminated wire-to-wire differences in gain due to differences in electronics, etc. Moreover, gain maps
had shown that the gain along a given sense wire was constant to within $\pm 2\%$, a variation considerably smaller than the effect which was measured here. Finally, a control run during which the strip source was completely shielded provided a measurement of the background, most of which came from cosmics. This background was not constant across the $^{55}\text{Fe}$ spectrum, but it was nevertheless small enough relative to the signal so that it could only shift the fitted mean of the $^{55}\text{Fe}$ spectrum by $0.2\%$ for all reasonable values of the mean. From these facts it could be assumed that the relative gain used in this analysis was a meaningful quantity.
By plotting the relative gain as a function of distance from the edge, and then repeating the whole procedure for various strip voltages $V_{\text{middle}}$ and $V_{\text{lower}}$, the optimal $V_{\text{middle}}$ and $V_{\text{lower}}$ were found. (N.B. In all cases $V_{\text{middle}}$ and $V_{\text{lower}}$ were held equal since unequal voltages did not provide any measurable advantage. Therefore, from now on $V_P$ will be used to refer to both $V_{\text{middle}}$ and $V_{\text{lower}}$.) To determine the optimal $V_P$ the average relative gain between 10 mm and 30 mm from the edge was calculated, and then the sum of the deviations $\delta$ from this average was found for relative gains between 1 mm and 10 mm from the edge. The point $\sim 0$ mm from the edge was not used because the main peak of the edge spectrum associated with it was not sufficiently distinct. For a set of measurements at $n$ values of $V_P$, the $n$ points obtained were fit to a straight line with $(x_i, y_i) = (V_{f_i}, \delta_i)$, and the $V_P$ corresponding to $\delta_i = 0$ was taken to be the optimal $V_P$. For example, to obtain the optimal voltage for the K-sector with $V_S = 1275$ and $V_P = 0$, measurements of the relative gain were made at several values of $V_P$: +50, +75, +100 and +125 Volts (see Fig. B.3). Then, the $\delta_i$ for each plot was found and a straight line was fit to the three points. The $V_P$ at which this line crosses $\delta_i = 0$ is taken as the optimal $V_P$.

### B.4 Results

Edge gain distortion measurements were made for two different K-sectors and one W-sector. The measurements were made on both edges of the sectors and at different positions along the edge to insure against spurious results caused by local variations in edge geometry. Two different sector types were used since in
the K-sector most of the wires make a 60° angle with the edge, whereas for the W (and M) sectors most of them make a 75° angle. In all three sectors there are edge regions where the angle is considerably different from those just mentioned, and measurements have shown that the optimal $V_P$ for these regions would be different. However, these regions comprise only a small percentage of the total number of wires in the sector, and since the priority in this measurement was naturally to minimize the gain distortion over the largest possible portion of the edge of the sector, only the $V_P$ 's for the main regions are reported here.

The optimal $V_P$ 's for several different $V_S$ and $V_F$ combinations (sense and
field wire voltages, respectively) were calculated. The "default" values for the TPC are $V_S = 1275$ and $V_F = 0$, but one can achieve higher sector stability [59] at the same gain by increasing $V_S$ and $V_F$ appropriately, thereby decreasing the electric field $E_F$ on the field wires while simultaneously maintaining the same gain levels [60]. For example, $V_S = 1407$ and $V_F = 295$ gives the same gain as the default values but better stability. Combinations of $V_S$ and $V_F$ corresponding to gains higher than the default were also checked. Tables B.1 and B.2 show the $V_S, V_F$ combinations tested, along with their corresponding $E_F$'s and calculated $V_{opt}$'s for the K and W sectors, respectively. These optimal voltages are a roughly linear function of $V_S$, as seen in Fig. B.4. Also tested was the effect of a $V_F$ set

<table>
<thead>
<tr>
<th>K Sector</th>
<th>$1 \times E_F$</th>
<th>$\frac{1}{3} \times E_F$</th>
<th>$0 \times E_F$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Gain</td>
<td>$V_S$</td>
<td>$V_F$</td>
<td>$V_{opt}$</td>
</tr>
<tr>
<td>1275</td>
<td>1275</td>
<td>0</td>
<td>119</td>
</tr>
</tbody>
</table>

Table B.1: Sense, field wire and optimal strip voltages for K sector for several field strengths (normalized to the field strength for $V_F = 0$) on the field wires. The left-most column gives the equivalent $1 \times E_F$ gain.

artificially high—perhaps useful if a sector becomes very unstable—and it produced gain distortions of about 20% within 2–3 mm of the edge which then became negligible 1 cm from the edge. Earlier work with the small chamber predicted that, for values of $V_S$ near 1275 the optimal $V_F$ would be somewhat bigger than
Field Strength on Field Wires

<table>
<thead>
<tr>
<th>W Sector</th>
<th>$1 \times E_F$</th>
<th>$\frac{1}{3} \times E_F$</th>
<th>$0 \times E_F$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Gain</td>
<td>$V_S$</td>
<td>$V_F$</td>
<td>$V_{opt}$</td>
</tr>
<tr>
<td>1275</td>
<td>1275</td>
<td>0</td>
<td>83</td>
</tr>
<tr>
<td>1300</td>
<td>1300</td>
<td>0</td>
<td>85</td>
</tr>
</tbody>
</table>

Table B.2: Sense, field wire and optimal strip voltages for W sector for several field strengths (normalized to the field strength for $V_F = 0$) on the field wires. The left-most column gives the equivalent $1 \times E_F$ gain.

the optimal $V_P$ actually found by this measurement. The different materials and geometry of the small chamber are probably responsible for this discrepancy.

B.5 Conclusion

Voltages for the middle and lower potential strips which minimize gain distortions have been found. They are a function of the sense and field wire voltages and the angle the wires make with the sector edge. Generally speaking, increasing the gain and/or $V_S, V_F$ requires an increase in $V_P$ to keep the gain flat near the edge.
Figure B.4: Optimal potential strip voltages as a function of sense wire voltage for K and W type sectors.
Bibliography


[10] The Higgs mechanism is described in detail in many sources, *e.g.* Quigg, *Gauge Theories*, 75.


[25] *Large Electron-Positron storage ring, Technical Notebook*, CERN Publica-
tions, European Laboratory for Particle Physics, 1211 Geneva 23, Switzer-

[26] D. Decamp *et. al.* (ALEPH Collaboration), *ALEPH: A Detector for
Electron-Positron Annihilations at LEP*, to be published in Nucl. Inst.
Meth.

[27] D. Decamp *et. al.* (ALEPH Collaboration), *Heavy Flavour Production in Z
Decays*, CERN-EP/90-54, accepted for publication by Phys. Lett. B.


[29] All hadronic processes were simulated with LUND 6.3 Parton Shower; M.
*et. al.*, Phys. Rev. **D37** (1988) 1. (This simulation has been shown to ac-
curately reproduce the data; D. Decamp *et. al.* (ALEPH Collaboration),

[30] For electron-pair production the BABAMC Monte Carlo was used; M.
and tau pair production the KORALZ program was used; S. Jadach, *et.
$2\gamma$ processes leading to $Z^0 \rightarrow l^+l^- (\gamma)$ final states the program DIAG36 was used; see F.A. Berends et. al. Nucl. Phys. B253 (1985) 441. For $2\gamma$ processes leading to hadronic final states the PHOT01 Monte Carlo, based on the work of J. A. M. Vermaseren, with later contributions from S. Kawabata, J. Olsson, H. Wriedt, and J. M. Nye, was used.


[32] T. Barklow, SLAC Report–315 (1987) 326. The approximations $m_i = m_j = 0$ and $E_j = 1$ GeV, where $E_j$ is the energy of the $j$th jet, are used.

[33] T. Sjöstrand, op. cit. A $d_{\text{join}}$ of 99999. is used.


[38] D. Decamp et. al. (ALEPH Collaboration), Measurement of Electroweak Parameters from Z Decays into Fermion Pairs, CERN-PPE/90-104; to be published in Z. Phys. C.


[41] B. Lynn and D. Kennedy, SLAC-PUB-4039.


[48] Ibid., The LUCLUS algorithm with a djoin of 2.8 is used.

[49] Ibid., The LUCLUS algorithm with a djoin of 99999. is used.


[51] (JADE Collaboration), Determination of Semi-Muonic Branching Ratios and Fragmentation Functions of Heavy Quarks in e+e− Annihilation at √s = 34.6 GeV, Z. Phys. C33 (1987) 339; J. Chrin, Upon the Determination of Heavy Quark Fragmentation Functions in e+e− Annihilations, DESY 87-040 (May 1987); S. Bethke, Recent Results on the Fragmentation

[52] Numerical Algorithms Group Ltd., NAG Central Office, Mayfield House, 256 Banbury Road, Oxford, OX2 7DE, UK. Routine E02AEF is used.


[54] NAG, op. cit. Routine E02DAF is used.


[59] Ron Settles, private communication.