SEARCHES FOR NEW PHYSICS

John Ellis and Felicitas Pauss

CERN, Geneva, Switzerland

To be published in
Proton–Antiproton Collider Physics
by World Scientific, Singapore, 1988

CERN TH.4992/88
March 1988
SEARCHES FOR NEW PHYSICS

John Ellis and Felicitas Pauss

CERN, Geneva, Switzerland

1. INTRODUCTION

The CERN p\bar{p} Collider has been the first accelerator to operate in a completely new energy domain, reaching centre-of-mass energies an order of magnitude larger than those previously available with the Intersecting Storage Rings (ISR) at CERN, or with the Positron–Electron Tandem Ring Accelerator (PETRA) at DESY and the Positron–Electron Project (PEP) at SLAC. Naturally there has been great interest in the searches for new physics in this virgin territory. Theorists have approached these searches from either or both of two rival points of view. Either they have had an a priori prejudice as to what new physics should be searched for, and what its signatures should be, or they have tried to interpret a posteriori some experimental observations. Whilst some experimentalists’ searches have been moulded by such a priori prejudices, many have emerged from systematic studies of all the measurable parameters of the events. The basic building-blocks of new physics in the 100 GeV energy domain are jets j, charged leptons \ell, photons \gamma, and missing transverse energy E_T. Therefore searches have been conducted in channels which are combinations of these elements. Table 1.1 lists the various combinations of j, \ell, \gamma, and E_T which have been explored by the UA1 and UA2 experiments. It also shows some of the main a priori theoretical prejudices which can be explored in each of these channels. Entries in Table 1.1 indicate the major searches which have been made, and which are reviewed here.

The layout of the rest of this paper is as follows. There are sections discussing each of the major prejudices: the Standard Model in Section 2; supersymmetry in Section 3; extra gauge degrees of freedom in Section 4; composite models in Section 5; and other possibilities in Section 6. Each of these sections contains a description of the motivations and characteristics of the new physics to be searched for, followed by a review of the searches made up to now at the CERN p\bar{p} Collider. Finally, Section 7 summarizes the lessons to be learnt so far from searches for new physics at the CERN p\bar{p} Collider, and previews some of the prospects for the next rounds of collider searches at CERN and FNAL.
Table 1.1
Possible new physics and signatures

<table>
<thead>
<tr>
<th></th>
<th>Standard Model</th>
<th>Supersymmetry</th>
<th>Extra gauge bosons</th>
<th>Composite models etc.</th>
</tr>
</thead>
<tbody>
<tr>
<td>j-j</td>
<td></td>
<td></td>
<td>Z' → q̅q</td>
<td>Λ_{contact}</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>W'</td>
<td>q'</td>
</tr>
<tr>
<td>j-ℓ</td>
<td>t, b'</td>
<td></td>
<td></td>
<td>Leptoquark</td>
</tr>
<tr>
<td>j E_T</td>
<td>N_r</td>
<td>ĝ, ̅ĝ</td>
<td></td>
<td>X → ZW</td>
</tr>
<tr>
<td></td>
<td>W → L_{ν_L}</td>
<td></td>
<td>ZZ</td>
<td></td>
</tr>
<tr>
<td>ℓℓ</td>
<td></td>
<td></td>
<td>Z' → ℓℓ</td>
<td>Leptoquark</td>
</tr>
<tr>
<td>ℓ E_T</td>
<td>W → L_{ν_L}</td>
<td>W → ℓ± ̅ν</td>
<td>W' → ℓν</td>
<td></td>
</tr>
<tr>
<td>jℓ E_T</td>
<td>t, b'</td>
<td>̃W, ̃Z</td>
<td>Z' → WW</td>
<td>X → WW</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>W' → ZW</td>
<td>ZZ</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td>Leptoquark</td>
</tr>
<tr>
<td>jℓℓ</td>
<td></td>
<td></td>
<td>W' → WZ</td>
<td>X → WZ</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>Z' → ZZ</td>
<td>ZZ</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td>Leptoquark</td>
</tr>
<tr>
<td>ℓℓ E_T</td>
<td>Z → L^+L^-</td>
<td>Z → ℓℓ</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>Z → ̃W̅W</td>
<td></td>
</tr>
<tr>
<td>γX</td>
<td></td>
<td></td>
<td>Z → ℓℓγ</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>W → ℓυγ</td>
<td></td>
</tr>
</tbody>
</table>

2. THE STANDARD MODEL

Within the minimal version of the Standard Model [2.1] with just three generations of quarks and leptons and just one Higgs doublet, there are just three more elementary particles waiting to be discovered, namely the ντ, the t-quark, and the neutral Higgs boson. Going beyond the minimal version, the Standard Model could easily accommodate more quark and lepton generations—implying more neutrinos, new charged heavy leptons L, and quarks, and/or more Higgs doublets—in turn implying more neutral Higgs bosons and also charged ones. Searches for more quarks at the CERN pp Collider are discussed elsewhere [2.2], as are constraints on the number N_ν of neutrino species inferred from the production rates for W± and Z^0 bosons [2.3]. In this section we concentrate on neutrino counting using events with large E_T, and on the searches for heavy leptons and the Higgs boson.

2
2.1 Neutrino Counting

Even though the \( \nu \) has not been observed directly, there are two indirect arguments for its existence. One is the agreement of the \( \tau \) lifetime and leptonic branching ratio with the Standard Model prediction \([2.4]\), and the other is the observation of \( W \to \tau \nu \) decay at a rate consistent with \( e-\mu-\tau \) universality of the weak charged-current coupling at \( Q^2 = m_W^2 \) \([2.3]\). Defining the ratio of coupling strengths \((g_1/g_2)^2 = \Gamma(W \to \ell_1 \nu \gamma) / \Gamma(W \to \ell_2 \nu \gamma)\), and combining the \( \sqrt{s} = 546 \text{ GeV} \) and 630 GeV \( W \to \ell \nu \) \((\ell = e, \mu, \tau)\) data samples, UA1 obtained \([2.5]\)

\[
g_\nu/g_\ell = 1.00 \pm 0.07 \pm 0.04, \quad g_\nu/g_e = 1.01 \pm 0.10 \pm 0.06,
\]

(2.1)

where the first error is statistical and the second one is due to systematic uncertainties. All the neutrinos are expected to be massless in the minimal version of the Standard Model, but could in principle acquire masses from extensions of it. Direct measurements of \( \tau \) decay tell us that \( m_\tau < 35 \text{ MeV}/c^2 \) \([2.6]\), and cosmological considerations using reasonable hypotheses about possible \( \nu \) decay modes then suggest that \( m_\nu \ll 100 \text{ eV}/c^2 \) \([2.7]\). Therefore any neutrino counting experiment with a kinematic range \( \geq 1 \text{ keV} \) should find \( N_\nu \geq 3 \).

An important historical role in bounding \( N_\nu \) has been played by cosmology through the successful confrontation of primordial nucleosynthesis calculations with observation \([2.8]\). The best fit is obtained with \( N_\nu = 3 \), but the upper bound on \( N_\nu \) depends on other inputs. If we assume \([2.9]\) (more conservatively \([2.10]\)) that the primordial \(^4\text{He} \) abundance \( Y(\text{He}) < 0.254 \) \((0.26)\), that the baryon-to-photon ratio \( N_b/N_\gamma \geq 3 \times 10^{-10} \) \((2 \times 10^{-10})\) as suggested by the abundance of \( D + \text{Li} \) \((\text{He}) \), and that the neutron half-life \( t_{\nu/2}(n) > 10.4 \) \((10.2)\) minutes, we find for the two assumptions

\[
N_\nu < 4.0 \ (5.2).
\]

(2.2)

(Keeping track of the decimal is not useful for counting conventional left-handed neutrinos, but it is useful when extending the analysis to other light neutral particle species, such as right-handed neutrinos.) More recently, astrophysics has provided an interesting upper bound on \( N_\nu \) thanks to the observation of neutrinos from the supernova SN 1987a \([2.11]\). The events seen are presumably due to \( \bar{\nu}_e \), and can be used to infer the energy output \( E_{\nu_e} \) of the supernova through \( \bar{\nu}_e \). As a first approximation, this can be multiplied by \( 2N_\nu \) to infer the total energy output through all neutrinos. This must be less than the binding energy of a neutron star, which is \( < 4 \times 10^{53} \text{ erg} \). Thus the data can be used to bound \( N_\nu \) \([2.12]\):

\[
N_\nu \leq 6.
\]

(2.3)

(We have rounded down to the nearest integer: any extension to right-handed neutrinos requires more input on the core of the neutron star.) The best fit to the supernova data is again with \( N_\nu = 3 \).
The best particle physics experimental limit, apart from those obtained at the CERN p̅p Collider, is provided by searches for the reaction \(e^+e^- \rightarrow \gamma + \text{nothing}\). At the present PEP and PETRA energies, the cross-section for this reaction is \(\propto (N_\nu + 4)\). A compilation of the world’s data (ASP, MAC, CELLO) gives \[N_\nu < 4.9 \quad (90\% \text{ CL}).\]  
(2.4)

This upper bound is increased to 7.3 if one assumes that there are at least three neutrino species \[N_\nu \geq 3\]. At the SLAC Linear Collider (SLC) and at the CERN Large Electron-Positron storage ring (LEP) the cross-section for this process just above the Z peak will be \(\propto N_\nu\), and it should be possible to obtain high enough statistics to measure \(N_\nu\) with a 10% error \[N_\nu \leq 5.9 \quad (90\% \text{ CL}).\]  
(2.5)

for \(m_\tau > 44\ \text{GeV}/c^2\) if one imposes the bound \(N_\nu \geq 3\). Details can be found in Ref. [2.16].

The other way to measure \(N_\nu\) at the CERN p̅p Collider is via the contribution to the missing-energy cross-section due to the process \(p̅p \rightarrow \text{gluon or quark} + (Z \rightarrow \nu\bar{\nu}) + X\), i.e. Z production at large \(p_T\). In this case the p̅p is transformed into large \(E_T\) since the neutrinos go undetected. We then observe the recoil gluon (or quark) which appears as a jet in the detector. The cross-section is \(\propto N_\nu\), and can be not only calculated using QCD but also calibrated by using gluon + \((W \rightarrow e\nu)\) events, i.e. W’s produced at large \(p_T\) [2.17]. However, there are significant backgrounds to high-p_T Z production: for example, from the production of heavy quarks (with subsequent semileptonic decay of the heavy quark, where the charged lepton goes undetected) and \(W \rightarrow \tau\nu\) (where the \(\tau\) decays hadronically), which produce genuine \(E_T\); and from instrumental effects due to mismeasurements of jet energies in QCD events with no \(\nu\) emission, which produce fake \(E_T\). Since missing-energy events play key roles in the searches for other new types of physics, e.g. heavy charged leptons L, squarks \(\tilde{q}\), and gluinos \(\tilde{g}\), we discuss here in more detail the UA1 1983 + 1984 + 1985 \(E_T\) event sample and possible background sources.

The aim was to define an event sample in which the background due to fluctuations in the detector response was small. The significance of the measured \(E_T\) was defined by

\[
N_\sigma = \frac{E_T}{0.7 \sqrt{\sum E_T}}, \quad E_T = \sum_i E_i^\parallel.
\]  
(2.6)
Fig. 2.1 Distribution of $N_{\pi}$ [Eq. (2.6)] for $> 3\sigma$ monojet events (only one jet with $E_T > 12$ GeV) shown as histogram, compared with all expected contributions (solid line) and jet fluctuation contributions only (shaded area) [2.19].

where all energies are in units of GeV, and $E_T^i$ is a vector in the direction of the calorimeter cell $i$ with magnitude $E_T^i$ and $\Sigma E_T$ is the scalar sum of transverse energy observed in all calorimeter cells. The background due to fluctuations of the detector response was evaluated with a Monte Carlo technique using real jet data [2.5]. Figure 2.1 shows the distribution of $N_{\pi}$ for events with $N_{\pi} > 3$ compared with all expected contributions (solid line) and jet fluctuation contributions only (shaded area). For $N_{\pi} < 4$ the observed event rate is dominated by contributions from jet fluctuations. Accordingly, an inclusive selection was made of events with isolated $E_T$ with $N_{\pi} > 4$, fulfilling the following conditions: i) $E_T > 15$ GeV and $N_{\pi} > 4$; ii) at least one jet with $E_T^j > 12$ GeV in $|\eta| < 2.5$ and a matching central detector (CD) track of $p_T > 1$ GeV/c; (iii) no e or $\mu$ candidates; (iv) no back-to-back jets, i.e. remove events with a second jet*) in an azimuthal angle region of $\pm 30^\circ$ opposite to the most energetic jet and around the $E_T$ direction; and (v) additional $E_T$ validation cuts were applied. A total of 56 events passed the selection cuts. The observed rate is in very good agreement with the contributions expected from Standard Model physics, as seen in the first column of Table 2.1. The calculations for all known physics processes were done with the ISAJET Monte Carlo program [2.18] including i) simulation of the full proton-antiproton collision to take into account the

*) Jet observed in the calorimeter with $E_T > 8$ GeV, or jet observed in the CD drift chambers with $p_T > 5$ GeV/c.
Table 2.1
Predicted rates for processes giving large $E_T$
using all event selection cuts [2.19]

<table>
<thead>
<tr>
<th>Process</th>
<th>Events (total)</th>
<th>Events with $L_T &lt; 0$</th>
<th>Events with $L_T &lt; 0$ and $E_T^1 &lt; 40$ GeV</th>
</tr>
</thead>
<tbody>
<tr>
<td>$W \rightarrow e\nu$</td>
<td>3.6</td>
<td>2.0</td>
<td>1.4</td>
</tr>
<tr>
<td>$\rightarrow \mu\nu$</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$\rightarrow \tau\nu \rightarrow$ leptons</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$W \rightarrow \tau\nu$</td>
<td>36.7</td>
<td>8.0</td>
<td>7.1</td>
</tr>
<tr>
<td>$\rightarrow \nu\bar{\nu} +$ hadrons</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$W \rightarrow c\bar{s}$</td>
<td>&lt; 0.1</td>
<td>&lt; 0.1</td>
<td>&lt; 0.1</td>
</tr>
<tr>
<td>$Z \rightarrow \tau^+\tau^-$</td>
<td>0.5</td>
<td>0.1</td>
<td>0.1</td>
</tr>
<tr>
<td>$Z \rightarrow \nu\bar{\nu}$</td>
<td>7.4</td>
<td>7.1</td>
<td>5.6</td>
</tr>
<tr>
<td>(3 neutrino species)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$Z \rightarrow c\bar{c}$ and $b\bar{b}$</td>
<td>&lt; 0.1</td>
<td>&lt; 0.1</td>
<td>&lt; 0.1</td>
</tr>
<tr>
<td>$c\bar{c}$ and $b\bar{b}$</td>
<td>0.2</td>
<td>0.2</td>
<td>0.2</td>
</tr>
<tr>
<td>(direct production)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Jet fluctuations</td>
<td>3.8</td>
<td>3.4</td>
<td>3.4</td>
</tr>
<tr>
<td>(fake missing energy)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>Total</td>
<td>52.2</td>
<td>20.8 ± 5.1 ± 1.0</td>
<td>17.8 ± 3.7 ± 1.0</td>
</tr>
</tbody>
</table>

The effects of the spectator particles, and ii) simulation of the UA1 detector including the hardware triggers. The accuracy of these Monte Carlo predictions has been checked by loosening the cuts, e.g. by relaxing to $N_\nu > 3$, and by removing the back-to-back cut for events with $N_\nu > 4$.

A scatter-plot of the 56 $E_T$ events remaining after the primary selection is shown in Fig. 2.2. The horizontal axis is the $E_T^1$ of the highest-$E_T$ jet, whilst the vertical axis is a measure of the $\tau$-likelihood $L_T$ of the event, defined by

$$L_T = \ln \left( P_T P_R P_N \right).$$ (2.7)
Fig. 2.2 Scatter plot of \( L_r \) [Eq. (2.7)] versus the transverse energy of the highest-\( E_T \) jet in the event. The projections for \( L_r \) and \( E_T \) are plotted for the data (histogram), for all predicted contributions (solid line), and for non-\( \tau \) contributions (shaded area). The different symbols indicate the charged multiplicity \( N \) [Eq. (2.10)] of the highest-\( E_T \) jet for the 56 events [2.19].

Here \( P_F \), \( P_R \), and \( P_N \) are the relative probabilities for values in the distributions of \( F \), \( R \), and \( N \):

— \( F \) is the fraction of the jet energy, as measured in the calorimetry

\[
F = \frac{\sum E_T \text{ in } \Delta R < 0.4}{\sum E_T \text{ in } \Delta R < 1},
\]  
(2.8)

where \( \Delta R = (\Delta \phi^2 + \Delta \eta^2)^{1/2} \); so that \( F \) is a measure of the collimation of the jet;

— \( R \) is the angular separation in \( \eta-\phi \) space,

\[
R = [(\phi_{T_1} - \phi_J)^2 + (\eta_{T_1} - \eta_J)^2]^{1/2},
\]  
(2.9)

where the subscript \( T_1 \) refers to the CD track with the largest \( p_T \), and \( J \) refers to the calorimeter jet axis; and

— \( N \) is the number of charged tracks in a cone of \( \Delta R = 0.4 \) about the calorimeter jet axis,

\[
N \equiv \text{No. of charged tracks (} p_T > 1 \text{ GeV}/c \).
\]  
(2.10)

Tau leptons produced in \( W \) decay have a large Lorentz boost, therefore the hadrons coming from the \( \tau \) decay form a narrow, high-\( p_T \) hadronic jet. In addition, \( \sim 99 \% \) of the hadronic
decays of the $\tau$ have charged multiplicities of one or three. These $\tau$ characteristics are combined in $L_\tau$ [Eq. (2.7)], which is used to select a low background sample of $W \rightarrow \tau \nu$ candidates. Events with $L_\tau > 0$ (32 events) can mainly be understood as $W \rightarrow \tau \nu$, decays with $\tau \rightarrow$ hadrons + $\nu$, and were used to extract the lepton universality result (2.1) [2.5]. The remaining 24 events with $L_\tau < 0$, which are less likely to be $W \rightarrow \tau \nu$, decays, have been used to search for new physics [2.19]. The second column of Table 2.1 shows Monte Carlo predictions for expected contributions to the $L_\tau < 0$ sample from known sources. The overall number agrees well with the observed number of 24 events in this category. Note that the contribution of jet + ($Z \rightarrow \nu \bar{\nu}$) events in Table 2.1 is estimated using the experimental distribution of large-$p_T$ W events [2.20] for normalization, and assuming just $N_\nu = 3$ neutrino flavours.

To establish an upper limit on $N_\nu$, the number of additional events expected with $L_\tau < 0$ and $E_\nu < 40$ GeV was calculated for each additional neutrino species: 1.8 event per species. Comparing this prediction with the numbers in the third column of Table 2.1, which already include three neutrino species, it was concluded [2.19] that

$$N_\nu \leq 10 \quad (90\% \, \text{CL}). \quad (2.11)$$

This result can alternatively be expressed as an upper limit on all $Z$ decays into any light neutral non-interacting particles (assuming a partial width of 180 MeV for each $Z \rightarrow \nu \bar{\nu}$ decay mode):

$$\Gamma(Z \rightarrow \text{nothing}) \leq 1.8 \, \text{GeV} \quad (90\% \, \text{CL}). \quad (2.12)$$

---

**Fig. 2.3** Summary of current limits on the number of neutrino species.

---

8
Note that in the calculation for the total expected background, no contribution from the top quark was included (i.e. $W \rightarrow t\bar{b}$, $Z \rightarrow t\bar{t}$, and $pp \rightarrow t\bar{t} + X$). Adding these contributions would improve the upper bound on the total number of neutrino species by an amount depending on $m_t$.

The limit (2.11) is not yet competitive with the most stringent of the other limits shown in Fig. 2.3, but it is more direct than all except the limits from $e^+e^- \rightarrow \gamma + \text{nothing}$. We can expect that the 'gluon-tagging' method for neutrino counting will be improved by data taken with the Antiproton Collector (ACOL) and the upgraded detectors, which should significantly reduce the statistical and systematic errors.

2.2 Fourth-Generation Charged Leptons

These are not included in the minimal three-generation version of the Standard Model, but we have no model-independent reason to exclude them theoretically. The pair-production of heavy leptons has been sought in $e^+e^-$ collisions at centre-of-mass energies $\sqrt{s} \leq 46.6\text{ GeV}$ with negative results, allowing us to conclude (2.13) that

$$m_L > 22.7\text{ GeV/c}^2 \quad (90\% \text{ CL}),$$

(2.13)

and improved upper limits are becoming available now from TRISTAN [2.21]. In principle, we could exclude indirectly the existence of any heavy lepton by determining that $N_L < 4$, but the bounds (2.2) to (2.5) and (2.11) on $N_L$ are not yet sufficiently precise. There is a weak upper bound on the possible mass of any heavy lepton from a global analysis of electroweak neutral- and charged-current data, which bounds radiative corrections due to a heavy lepton, telling us that

$$m_L < 290\text{ GeV/c}^2 \quad (90\% \text{ CL}),$$

(2.14)

if any such particle exists [2.22].

There are at least three ways of producing a heavy lepton at the CERN p$\bar{p}$ Collider: through $W \rightarrow L\nu_L$, through $\gamma^*, Z \rightarrow L^+L^-$, and through heavy-quark decays such as $t \rightarrow L\nu_Lb$. In practice, only the first channel has been studied [2.23], since it has a cross-section much larger than that of the second, and the t-quark has not yet been observed. A conventional sequential heavy lepton is expected to have the branching ratio

$$\Gamma(W \rightarrow L\nu_L)/\Gamma(W \rightarrow e\nu_e) = 1[1 - 1.5\rho + 0.5\rho^3], \quad \rho \equiv (m_L/m_W)^2,$$

(2.15)

for $m_L < m_W$, assuming $m_{\tau L} = 0$. Its total decay rate is expected to be

$$\Gamma(L \rightarrow \text{all}) = (G_F^2 m_L^3/192\pi^3)[3 + 6(1 + \alpha_s/\pi)],$$

(2.16)
where the first term corresponds to the almost equally probable $L \rightarrow \nu_L e\bar{e}_e$, $\nu_L \mu\bar{\nu}_\mu$, and $\nu_L \tau \bar{\nu}_\tau$ decay modes, and the second term corresponds to the hadronic decay modes $L \rightarrow \nu_L (u\bar{d}, c\bar{s})$. For $m_L = 40 \text{ GeV}/c^2$, the expression (2.16) gives a lifetime $\tau_L \approx 10^{-19} \text{ s}$, which is too short to leave a detectable track. The leptonic decay modes of the $L$ ($L \rightarrow \ell \nu_L \ell$, $\ell = e, \mu, \tau$) would be difficult to disentangle from the $W \rightarrow e\bar{e}_e, \mu\bar{\nu}_\mu$ and $W \rightarrow \tau \bar{\nu}_\tau (\tau \rightarrow \ell \nu_\ell \nu_\ell)$ backgrounds, which have larger rates. Accordingly the search performed by UA1 has focused on the $L \rightarrow \nu_L q\bar{q}$ decays [2.19] with a total branching ratio of 67% and with an expected event topology of jet(s) + $E_T$.

This decay pattern gives events which are qualitatively similar to $W \rightarrow \tau \nu_\tau (\tau \rightarrow \text{hadrons} + \nu_\tau)$ except for the phase-space factor (2.15) and the typically larger invariant mass of the $q\bar{q}$ system. The fragmentation into hadrons has been modelled using a modified version of ISAJET where spin effects have been included [2.24]. Since heavy leptons come from $W$ decays, such $W \rightarrow L\bar{\nu}_L$ events would mostly populate the $E_T^L < 40 \text{ GeV}$ region of Fig. 2.2. Also, the wider jets fall mainly into the $L_\tau < 0$ region, as can be seen in Fig. 2.4 for heavy-lepton Monte Carlo events ($m_L = 55 \text{ GeV}/c^2$). Figure 2.5 shows the number of events expected from $W \rightarrow L\bar{\nu}_L \rightarrow q\bar{q}\nu_L\bar{\nu}_L$ in the $L_\tau < 0$, $E_T^L < 40 \text{ GeV}$ domain of Fig. 2.4, as a function of $m_L$.

The net acceptance for a heavy lepton of mass $40 \text{ GeV}/c^2$, decaying semihadronically, is 6.2%. This low efficiency comes mainly from the trigger efficiency and the $N_T > 4$ requirement. With the jet definition used, most of the heavy-lepton decays are predicted to give monojet events. The fraction of dijet events ($E_T^L > 12 \text{ GeV}$) ranges from $\sim 18\%$ ($m_L = 25 \text{ GeV}/c^2$) to $\sim 27\%$ ($m_L = 65 \text{ GeV}/c^2$).

![Fig. 2.4 Scatter plot of $L_\tau$ versus $E_T^L$ of the highest-$E_T$ jet for the heavy-lepton Monte Carlo events with $m_L = 55 \text{ GeV}/c^2$. The symbols indicate the charged track multiplicity as defined in Eq. (2.10).](image)

10
Fig. 2.5 The rate of heavy-lepton events passing the isolated $4\sigma E_T$ selection and the cuts $L_T < 0$ and $E_T^L < 40$ GeV as a function of the heavy-lepton mass (solid points). Also indicated is the lower limit on $m_L$ from $e^+e^-$ experiments.

Using the calculated heavy-lepton contributions and including the contribution of one additional neutrino $\nu_L$ to missing-energy events from jet + ($Z \to \nu_L\bar{\nu}_L$), it was found [2.19] that

$$m_L > 41 \text{ GeV/c}^2 \quad (90\% \text{ CL}).$$

(2.17)

Note that in computing this limit, no $t$-quark contributions were added to the background, and $L \to t\bar{b}f$ decays were not included in the predicted event rates. These contributions, if included, would increase the limit.

Studies have shown that the fraction of dijet + $E_T$ events increases as the mass of the heavy lepton increases [2.24]. These events are expected to have a more distinctive signature: two high-$p_T$ jets balanced by $E_T$ in a 'Mercedes'-type configuration. This event topology is not expected to be dominant in dijet events coming from the standard physics processes. Therefore the improved statistics and systematics of data accumulated with ACOL (including better calorimeter granularity of the upgraded UA1 detector) should make it possible to search for a signal for high-mass heavy leptons ($m_L \geq 45$ GeV/c$^2$) in the dijet event topology also. In the absence of a signal, the limits are expected to be improved considerably—in principle up to the phase-space limit of the $W$ decay (2.15).

2.3 The Higgs Boson

The presence of at least one physical neutral Higgs boson $H^0$ is essential for the renormalizability of the Standard Model. In the minimal version of the Standard Model with just one Higgs doublet, there is a single $H^0$ with well-defined couplings:
\[ g_{H\ell} = g_{\mu}/2m_{\ell}, \]  
\[ g_{HWW^+} = g_{m_{\ell}}, \]  
\[ g_{HZZ} = g_{m_{\ell}}, \]  
(2.18a)  
(2.18b)  
(2.18c)

where \( G_{\mu}/\sqrt{2} = g_{\mu}/8m_{\ell} \). Since its couplings increase with the mass of the other particles, we expect the Higgs to decay into the heaviest available particles. For example, a Higgs with mass between 11 GeV/c\(^2\) and 2m\(_{\ell}\) or 2m\(_{W}\) has

\[ \text{BR}(H \rightarrow e^+e^-: \mu^+\mu^-: \tau^+\tau^-: c\bar{c}: b\bar{b}) = 1: m_{\ell}^2/m_{\nu}^2: m_{\mu}^2/m_{\nu}^2: 3m_{\tau}^2/m_{\nu}^2: 3m_{c}^2/m_{\nu}^2. \]

Whilst the couplings of this minimal Higgs boson are known, its mass is not. Radiative corrections in the minimal version of the Standard Model suggest [2.25]

\[ m_H \gtrsim (O(\alpha_m^2))^{1/2} \approx \text{a few GeV}/c^2, \]  
(2.19)

but this can be evaded if \( m_{\nu} \approx 100 \text{ GeV}/c^2 \). Imposing tree-level unitarity on WW scattering amplitudes suggests [2.26] that

\[ m_H \lesssim (O(\alpha_m^2))^{1/2} \approx 1 \text{ TeV}/c^2. \]  
(2.20)

However, tree-level unitarity does not necessarily hold if the Higgs sector is strongly coupled, and although it has been argued that some observable effect should show up in the 1 TeV/c\(^2\) mass range even in this case, it would surely not take the form of an identifiable narrow particle state. Phenomenologically, one knows from the absence of Higgs-boson effects in low-energy neutron–nucleus scattering, nuclear decays, and the spectra of muonic X-rays, that [2.27]

\[ m_H \gtrsim 15 \text{ MeV}/c^2. \]  
(2.21)

With a view to searching for heavier Higgs bosons, estimates have been made of the branching ratio for \( \eta' \rightarrow \eta + H \) [2.28] and \( \psi' \rightarrow J/\psi + H \) [2.29], and of the flavour-changing vertices for \( s \rightarrow d + H \) and \( b \rightarrow s + H \) [2.30] which control the branching ratios for \( K \rightarrow \pi + H \) and \( B \rightarrow K + H \), respectively. Of these, the branching ratio for \( \eta' \rightarrow \eta + H \) is difficult to estimate reliably, and the experimental upper limit [2.31] on \( \psi' \rightarrow J/\psi + H \) is not sensitive enough to be interesting. The flavour-changing vertices for \( s \rightarrow d + H \) and \( b \rightarrow s + H \) can, in principle, be calculated reliably, but in practice the available calculations [2.30] disagree amongst themselves and with a previous general theorem [2.32] on the form of such vertices.
which should be satisfied. Some of the existing calculations have been used to argue that unsuccessful searches for \( K^+ \rightarrow \pi^+ + H \) and \( B \rightarrow K + H \) [2.33] exclude certain ranges of Higgs masses [2.34, 2.35]. We believe these conclusions to be premature, and await the results of a new calculation of the flavour-changing Higgs vertices which is now under way [2.36]. However, we note that any vertex of the general form previously derived would probably yield branching ratios for \( K^+ \rightarrow \pi^+ + H \) and \( B \rightarrow K + H \) that are below present experimental sensitivities.

Although searches for a Higgs boson in \( J/\psi \rightarrow \gamma + X \) have not yet reached sufficient sensitivity to rule out \( m_H \leqslant m_{J/\psi} \), searches for a Higgs boson in a combined sample of \( \Upsilon \rightarrow \gamma + X \) and \( \Upsilon'' \rightarrow \gamma + X \) decays [2.37] have now reached sufficient sensitivity to rule out a Standard Model Higgs in the range

\[
0.6 \leqslant m_H \leqslant 3.9 \text{ GeV}/c^2,
\]

(2.22)
even after QCD radiative corrections [2.38] to the branching ratio are included. This is the only range of Higgs-boson masses above 15 MeV that seems to be reliably excluded.

Two ways to search for a Higgs boson at the CERN \( p\bar{p} \) Collider have been considered. They are \( p\bar{p} \rightarrow (W^+ \rightarrow (W + H)) + X \) or \( W \rightarrow (W + H)) + X \), and \( p\bar{p} \rightarrow H + X \). The former exploits a direct-channel \( W \) pole and the relatively large coupling (2.18b). However, even for light Higgses in the Standard Model the number of such events is expected to be \( \lesssim 2 \times 10^{-3} \) [2.39]. Since only about 400 \( W \) events have been detected in UA1 and about 250 \( W \) events in UA2, neither experiment yet has sensitivity to the \((W + H)\) process. Even with more \( W \) events, it would not be easy to pick out any accompanying Higgs from the conventional hadronic background, although if \( m_H > 2m_b \) this would be less difficult with a microvertex detector if it allows the b-quark to be tagged. The dominant mechanism for unaccompanied Higgs production at the CERN \( p\bar{p} \) Collider is expected to be gluon–gluon fusion: \( gg \rightarrow H \) via a t-quark loop. The cross-section for this reaction to leading order in \( \alpha_s \) is [2.40]

\[
\sigma \ [p\bar{p} \rightarrow (gg \rightarrow H) + X] = \left( \frac{\pi}{32} \right) (\alpha_s/\pi)^2 (G_F/\sqrt{2}) (N^2/9) \tau \int d\tau \ \Gamma(\sqrt{\tau} e^\gamma) \ \Gamma(\sqrt{\tau} e^{-\gamma}),
\]

(2.23)

where \( \tau = m_H^2/s \), \( G(\xi) \) is the gluon distribution function, and \( N = 1 \) if the t-quark is the only quark much heavier than \( m_H \). Unfortunately, the expression (2.23) gives unobservably low rates for \( H \) production in the absence of a clear decay signature, whatever the values of \( m_H \).

Some excitement was caused by the final analysis of events in the \( W + \) jet(s) data sample of UA1 [2.20]. Two events were observed at very large transverse momentum of the \( W \) (\( p_T^W = 100 \text{ GeV}/c \)) in which the \( W \) candidate recoils against an energetic dijet system. The mass of the jet-jet system is compatible with the \( W \) mass. These events could be interpreted as \( WW \), \( WZ \) or \( ZZ \) pair-production, and one might ask whether they could be due to Higgs production and decay into \( WW \) and/or \( ZZ \). However, the cross-section for producing a
Higgs boson at the CERN pp Collider is very small \(2.40\). As an example, UA1 estimated the expected number of events for \(H^0 \rightarrow W^+W^-\) passing all selection cuts and being consistent with the parameters of the observed events, to be \(\leq 2 \times 10^{-3}\). Therefore, neither in the present Collider data sample nor in the future data set accumulated with ACOL does one expect an observable event rate from Higgs boson decay into \(W^+W^-\) or ZZ.

3. SUPERSYMMETRY

There are many reasons for postulating a supersymmetric extension of the Standard Model, starting with the ‘fundamental’ ones of elegance or of utility in constructing a unified Theory of Everything (TOE). However, the only motivation for having light supersymmetric particles in an accessible mass range is as a solution to the naturalness or hierarchy problem \([3.1]\). This has its roots in the fact that radiative corrections to the masses of elementary scalar bosons, such as the Higgs boson of the Standard Model, are quadratically divergent,

\[
\delta m^2_H = O(\alpha/\pi) \Delta^2, \tag{3.1}
\]

where \(\Delta\) is a loop momentum cut-off. It can be seen from Eq. (3.1) that in order for the physical mass of the Higgs to be natural, i.e. \(\delta m^2_H \ll m^2_H\), where \(m^2_H = m^2_H\), \(\Delta\) should be \(\leq 1\) TeV. The supersymmetric solution to this naturalness problem starts from the observation that the quadratically divergent boson and fermion loop corrections (3.1) to \(m^2_H\) have opposite signs, and that if bosons \(B\) and fermions \(F\) occur in pairs with identical couplings as in supersymmetry, these quadratically divergent corrections will cancel, leaving behind a residual

\[
\delta m^2_H = O(\alpha/\pi) |m^2_B - m^2_F|. \tag{3.2}
\]

Therefore the role of the cut-off \(\Delta\) in Eq. (3.1) is taken over by the boson–fermion mass difference. In this approach

\[
|m^2_B - m^2_F| \leq 1\ \text{TeV}/c^2, \tag{3.3}
\]

hence one expects the supersymmetric partners of the known particles listed in Table 3.1 to have masses \(\leq 1\) TeV/c^2.

The gauge and Yukawa couplings of the known particles have the following counterparts for their supersymmetric partners:

\[
g_1\gamma_\mu V^\mu \rightarrow \sqrt{2}g_1\gamma_\mu V^\mu + \text{h.c.,} \quad g_1^2 f^2 \bar{f}V^\mu, \quad (g_1^2/2)|f^2|^2 \tag{3.4a}
\]
\[
\lambda^2 \bar{H} H + \lambda^2 \bar{H} H + \ldots, \quad \lambda^2 (|\bar{H} H|^2 + |F H|^2). \tag{3.4b}
\]

It is easy to see that these interactions all conserve multiplicatively a quantum number called R-parity which is defined to be +1 for Standard Model particles and −1 for all their supersymmetric partners [3.2]. It is possible to relate R-parity to other quantum numbers conserved in the Standard Model:

\[
R = (-1)^{B + 3L + 2S}. \tag{3.5}
\]

**Table 3.1**

Spectrum of SUSY particles

<table>
<thead>
<tr>
<th>Particle</th>
<th>Spin</th>
<th>Sparticle</th>
<th>Spin</th>
</tr>
</thead>
<tbody>
<tr>
<td>quark $q_{l,R}$</td>
<td>$\frac{1}{2}$</td>
<td>squark $\bar{q}_{l,R}$</td>
<td>$0$</td>
</tr>
<tr>
<td>lepton $\ell_{l,R}$</td>
<td>$\frac{1}{2}$</td>
<td>slepton $\bar{\ell}_{l,R}$</td>
<td>$0$</td>
</tr>
<tr>
<td>photon $\gamma$</td>
<td>$1$</td>
<td>photino $\bar{\gamma}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>gluon $g$</td>
<td>$1$</td>
<td>gluino $\bar{g}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>W</td>
<td>$1$</td>
<td>wino $\bar{W}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>Z</td>
<td>$1$</td>
<td>zino $\bar{Z}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>Higgs $H$</td>
<td>$0$</td>
<td>shiggs $\bar{H}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>graviton $G$</td>
<td>$2$</td>
<td>gravitino $\bar{G}$</td>
<td>$\frac{3}{2}$</td>
</tr>
</tbody>
</table>

The conservation of R-parity has three important phenomenological consequences.

i) Sparticles are always produced in pairs, e.g. $e^+e^- \rightarrow \mu^+\mu^-$, $p\bar{p} \rightarrow (\bar{g}g$ or $\bar{q}q$ or $\bar{q}g) + X$. This property can also be seen directly from the forms of the interactions (3.4).

ii) Every sparticle decays into another sparticle, e.g. $\bar{e} \rightarrow e + \bar{\gamma}$, $\bar{g} \rightarrow q\bar{q}$ or $q\bar{q}\bar{\gamma}$, $\bar{q} \rightarrow q + \bar{g}$ or $q + \bar{\gamma}$.

iii) The lightest supersymmetric particle (LSP) is absolutely stable, since it has no legal decay mode.

Since the LSP is absolutely stable, it should be present in large numbers in the Universe today as a cosmological relic from the Big Bang. If the relic LSPs had strong and/or electromagnetic interactions, they would have bound with conventional nuclei to form exotic isotopes. These have not been seen [3.3] even at a level far below the standard predictions for LSP abundances [3.4]. Therefore we conclude that the LSP has no strong or electromagnetic interactions; thus it can escape from experimental detection in much the same way as a conventional neutrino. Just as $E_T$ was used as a signature in the search for $W \rightarrow ev, \mu\nu, \tau\nu$, so $E_T$ could also be a signature for sparticle pair-production.
In Table 3.1 there are many neutral weakly interacting sparticles which are candidates for being the LSP. They include the sneutrinos $\tilde{\nu}_e$, $\tilde{\nu}_\mu$, $\tilde{\nu}_\tau$, the shiggs $\tilde{H}$, photino $\tilde{\chi}$, and zino $\tilde{Z}$ of spin $1/2$, and the gravitino $\tilde{G}$ of spin $3/2$. In fashionable models the sneutrinos $\tilde{\nu}$ and gravitino $\tilde{G}$ are usually heavier than the lightest spin-$1/2$ sparticle, and unsuccessful dark-matter searches squeeze the $\tilde{\nu}$ possibilities still further. Hence the LSP seems most likely to be some $\tilde{H}/\tilde{\chi}/\tilde{Z}$ mixture [3.5]. Model studies easily produce examples where the LSP is an almost pure $\tilde{H}$ or $\tilde{\gamma}$ state. The cosmological relic density of an $\tilde{H}$ LSP tends to be too high unless $m_{\tilde{H}} > m_\nu$ or $m_r$. Thus there is a tendency to assume that the LSP is (at least approximately) a $\tilde{\gamma}$, although this is by no means ironclad. Searches for sparticle production have indeed generally assumed that the LSP is the $\tilde{\gamma}$ and that $m_\gamma$ is negligible compared with the other sparticle masses, which will also be our working hypothesis here. In most cases, the limits are essentially unchanged for $m_\gamma$ up to half the decaying sparticle masses [3.6].

### 3.1 Squarks and Gluinos

Since these sparticles have strong interactions, they have the largest cross-sections at hadron–hadron colliders. Unsuccessful searches in $e^+e^-$ collisions [3.7] tell us that

$$m_{\tilde{q}} \gtrsim 21.5 \text{ GeV/c}^2 \quad (90\% \text{ CL}). \quad (3.6)$$

It is normally assumed [3.8, 3.9] that the supersymmetric partners of the left- and right-handed quarks have essentially equal masses,

$$m_{\tilde{q}_L} = m_{\tilde{q}_R}, \quad (3.7)$$

and that the masses of the $\tilde{u}$, $\tilde{d}$, $\tilde{s}$, $\tilde{c}$, and $\tilde{b}$ squarks are also essentially equal,

$$m_{\tilde{t}} = m_{\tilde{b}} = m_{\tilde{t}} = m_{\tilde{c}} = m_{\tilde{b}}. \quad (3.8)$$

Neither of these assumptions is above suspicion, since unequal vacuum expectation values (v.e.v.’s) for the two Higgs doublets in the minimal supersymmetric extension of the Standard Model would split the squark masses. However, if these v.e.v.’s were equal, one would expect $|m_{\tilde{q}_1} - m_{\tilde{q}_2}| = |m_{\tilde{q}_1} - m_{\tilde{q}_2}| \ll m_{\tilde{q}}$, and this assumption of degeneracy would be justified. When $m_{\tilde{q}} > m_\gamma$, we expect $\tilde{q} \rightarrow q\tilde{\gamma}$ to dominate, whereas when $m_{\tilde{q}} < m_\gamma$, a large branching ratio for $\tilde{q} \rightarrow q\tilde{\gamma}$ decay is expected, although $\tilde{q} \rightarrow q\tilde{W}$ or $q\tilde{Z}$ decay could also be important. The implications of non-degenerate squark masses and of $\tilde{q} \rightarrow q\tilde{W}$, $\tilde{Z}$ decays for hadron–hadron searches are discussed elsewhere [3.10, 3.11]. Here we assume both degenerate squark masses and the dominance of $\tilde{q} \rightarrow q\tilde{\gamma}$ decays. We mainly consider the case where the $\tilde{\gamma}$ is stable, but the case where $\tilde{\gamma} \rightarrow \gamma + \tilde{H}$ will also be discussed.
Fig. 3.1 Missing transverse energy spectrum from $p\bar{p} \rightarrow \tilde{g}\tilde{g} + X$ for different values of $m_{\tilde{g}}$ at $\sqrt{s} = 630$ GeV.

Fig. 3.2 Cross-section for $p\bar{p} \rightarrow \tilde{q}\tilde{q} + X$ production as a function of $m_{\tilde{q}}$ and a fixed gluino mass of 4 GeV/c^2.

The lower limits on $m_{\tilde{g}}$ from fixed-target hadron-hadron collisions, from $\Upsilon$ decays, and from bottomonium $^3P_1_l(b\bar{b}) \rightarrow g + \tilde{g}\tilde{g}$ searches are in the range [3.7, 3.12]

$$m_{\tilde{g}} \geq (2\text{ to } 5) \text{ GeV/c}^2,$$  \hspace{1cm} (3.9)

depending on the squark mass. When $m_{\tilde{g}} > m_{\tilde{q}}$, we expect $\tilde{g} \rightarrow \tilde{q}\tilde{q}$ or $q\bar{q}$ decays to dominate, whereas when $m_{\tilde{g}} < m_{\tilde{q}}$, we expect $\tilde{g} \rightarrow q\bar{q}\Upsilon$ to dominate. Thus gluino decays always give rise to the final state $q\bar{q}\Upsilon$, whilst squark decays produce the dominant final states $q\Upsilon$ (if $m_{\tilde{g}} > m_{\tilde{q}}$) or $q\bar{q}\Upsilon$ (if $m_{\tilde{g}} > m_{\tilde{q}}$). Therefore, we expect multi-jets and large $E_T$ as an event signature.

When $m_{\tilde{g}} \ll m_{\tilde{q}}$, the dominant production process is $p\bar{p} \rightarrow \tilde{g}\tilde{g} + X$, which produces the $E_T$ spectrum shown in Fig. 3.1. However, the associated production process $p\bar{p} \rightarrow \tilde{g}\tilde{g} + X$ can also be important, although its cross-section falls rapidly with increasing $m_{\tilde{q}}$, as seen in Fig. 3.2. The relative importance of the different $\tilde{g}\tilde{g}$, $\tilde{q}\tilde{q}$, and $q\bar{q}$ production processes is shown in Table 3.2 for different choices of $m_{\tilde{g}}$ and $m_{\tilde{q}}$.

When $m_{\tilde{g}} < m_{\tilde{q}}$, $\tilde{q}\tilde{q}$ pair-production dominates the production of strongly interacting particles, whereas $\tilde{g}\tilde{g}$ dominates if $m_{\tilde{q}} > m_{\tilde{g}}$, as illustrated in Fig. 3.3. Production of associated $\tilde{q}\tilde{q}$ pairs can also be important if $m_{\tilde{q}} \approx m_{\tilde{g}}$, as also seen in Table 3.2, and $\tilde{g}\Upsilon$ and $\tilde{q}\Upsilon$ final states are not completely negligible\footnote{For instance, direct photino production $p\bar{p} \rightarrow \tilde{\chi}\tilde{q}$ contributes $\approx 3\%$ for $m_{\tilde{g}} = 90$ GeV/c^2, $m_{\tilde{q}} = 50$ GeV/c^2, increasing to $\approx 20\%$ of the total cross-section as $m_{\tilde{g}} \rightarrow \infty$.}. All these processes should be included in a
Table 3.2
Relative cross-section contributions from gluino and squark production
($p\bar{p} \rightarrow \tilde{g}\tilde{g}, \tilde{g}\tilde{q}, \tilde{q}\tilde{q} + X$) at $\sqrt{s} = 630$ GeV
for selected values of the gluino and squark masses

<table>
<thead>
<tr>
<th>Mass (GeV/c$^2$)</th>
<th>Fraction (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$m_{\tilde{g}}$</td>
<td>$m_{\tilde{q}}$</td>
</tr>
<tr>
<td>120</td>
<td>50</td>
</tr>
<tr>
<td>50</td>
<td>120</td>
</tr>
<tr>
<td>70</td>
<td>90</td>
</tr>
</tbody>
</table>

Fig. 3.3 Schematic picture of supersymmetric processes dominating different areas of the $m_{\tilde{g}}$-$m_{\tilde{q}}$ plane.

complete analysis of squark and gluino production at a hadron-hadron collider, also taking into account QCD radiative corrections if possible, and certainly in the analysis of light $\tilde{g}\tilde{g}$ pair-production, where they can be important. The results of analyses are generally published as allowed domains in the ($m_{\tilde{g}}, m_{\tilde{q}}$) plane.

We start by describing the squark and gluino search by the UA1 Collaboration [3.13]. A special selection was made to optimize the signal-to-background ratio for heavy $\tilde{g}$ and $\tilde{q}$ and for light $\tilde{g}$, taking into account the complexity of different topologies depending on the $\tilde{g}$ and $\tilde{q}$ mass combinations. The standard missing-energy event selection described in subsection 2.2
was used ($E_T > 15 \text{ GeV}$ and $N_e > 4$), and the subsample with $L_r < 0$ was selected to reduce contamination from $W \rightarrow \tau \nu$, ($\tau \rightarrow \text{hadrons} + \nu$). Two changes were made in the event selection: i) instead of at least one jet with $E_T^j > 12 \text{ GeV}$, at least two jets were required; ii) instead of the stringent cut on the isolation of the $E_T$ vector, the difference $\Delta\phi$ in azimuthal angle between the two highest-$E_T$ jets was required to satisfy $\Delta\phi < 140^\circ$. A total of four events passed all the above selection criteria. The expected contributions from Standard Model processes and from jet fluctuation background were evaluated using ISAJET and the jet fluctuation Monte Carlo as described in subsection 2.2. The predictions for the numbers of events from different sources are shown in Table 3.3. The total expected number of $5.2 \pm 1.9 \pm 1.0$ events should be compared with the four events observed. Figure 3.4 shows the

<table>
<thead>
<tr>
<th>Process</th>
<th>No. of events</th>
</tr>
</thead>
<tbody>
<tr>
<td>W/Z decays</td>
<td>3.0</td>
</tr>
<tr>
<td>$b\bar{b}$/c$\bar{c}$</td>
<td>2.0</td>
</tr>
<tr>
<td>Jet fluctuations</td>
<td>0.2</td>
</tr>
<tr>
<td><strong>Total</strong></td>
<td>$5.2 \pm 1.9 \pm 1.0$</td>
</tr>
</tbody>
</table>

*) The first error is statistical, the second is systematic.

![Graph](image)

Fig. 3.4 Distribution of $\Delta\phi$, the azimuthal angle between the two highest-$E_T$ jets in the event: for data (histogram), for the expectation from conventional processes plus background (dashed curve), and for squark and gluino production (dot-dashed curve), using $m_{\tilde{q}} = 60 \text{ GeV/c}^2$ and $m_{\tilde{g}} = 70 \text{ GeV/c}^2$. The supersymmetric particle production was multiplied by a factor of 10 [3.13].
Table 3.4
Production cross-sections, experimental acceptances, and predicted event rates from squark and gluino production (p\bar{p} \rightarrow \tilde{g}\tilde{g}, \tilde{q}\tilde{q}, \tilde{g}\gamma, \tilde{q}\gamma + X) for selected values of the squark and gluino mass. The errors quoted are the statistical errors resulting from the Monte Carlo generation [3.13].

<table>
<thead>
<tr>
<th>m_{\tilde{g}} (GeV/c^2)</th>
<th>m_{\tilde{q}} (GeV/c^2)</th>
<th>Cross-section (nb)</th>
<th>Acceptance (%)</th>
<th>Events</th>
</tr>
</thead>
<tbody>
<tr>
<td>30</td>
<td>200</td>
<td>20.20</td>
<td>0.2</td>
<td>24.0 ± 5.6</td>
</tr>
<tr>
<td>50</td>
<td>120</td>
<td>0.71</td>
<td>1.3</td>
<td>6.8 ± 0.8</td>
</tr>
<tr>
<td>50</td>
<td>90</td>
<td>0.78</td>
<td>1.4</td>
<td>7.7 ± 0.9</td>
</tr>
<tr>
<td>60</td>
<td>70</td>
<td>0.49</td>
<td>3.2</td>
<td>11.2 ± 0.8</td>
</tr>
<tr>
<td>70</td>
<td>90</td>
<td>0.11</td>
<td>2.5</td>
<td>2.0 ± 0.2</td>
</tr>
<tr>
<td>200</td>
<td>30</td>
<td>4.35</td>
<td>0.6</td>
<td>17.0 ± 2.2</td>
</tr>
<tr>
<td>120</td>
<td>50</td>
<td>0.44</td>
<td>3.8</td>
<td>12.0 ± 0.8</td>
</tr>
<tr>
<td>90</td>
<td>50</td>
<td>0.64</td>
<td>3.8</td>
<td>17.5 ± 1.2</td>
</tr>
<tr>
<td>70</td>
<td>60</td>
<td>0.49</td>
<td>6.0</td>
<td>20.9 ± 1.2</td>
</tr>
<tr>
<td>130</td>
<td>70</td>
<td>0.07</td>
<td>6.6</td>
<td>3.3 ± 0.2</td>
</tr>
<tr>
<td>90</td>
<td>70</td>
<td>0.13</td>
<td>7.0</td>
<td>6.5 ± 0.3</td>
</tr>
<tr>
<td>70</td>
<td>70</td>
<td>0.27</td>
<td>5.3</td>
<td>10.2 ± 0.8</td>
</tr>
</tbody>
</table>

\( \Delta \phi \) distribution for the large-\( E_T \) multijet sample before applying the cut \( \Delta \phi < 140^\circ \). As can be seen, most events with two jets are almost back-to-back in azimuth, as are the background contributions (i.e. heavy flavour and jet fluctuations), whilst the angular distribution expected from a typical supersymmetric scenario is quite flat. The contributions to the \( \Delta \phi < 140^\circ \) multijet sample from squark and gluino production were again evaluated using the ISAJET Monte Carlo program, with full simulation of the UA1 detector including hardware triggers. Events were generated for a wide range of squark and gluino masses. Examples of predicted event rates and efficiencies are given in Table 3.4.

The selection efficiency is relatively high for large values of \( m_{\tilde{g}} \) (\( \epsilon = 3\% \) for \( m_{\tilde{g}} = 70 \text{ GeV/c}^2 \), \( m_{\tilde{q}} \rightarrow \infty \)) but quite small for small values of \( m_{\tilde{g}} \) (\( \epsilon = 0.01\% \) for \( m_{\tilde{g}} = 20 \text{ GeV/c}^2 \), \( m_{\tilde{q}} \rightarrow \infty \)), if only the lowest-order 2-to-2 production processes \( gg \rightarrow \tilde{g}\tilde{g} \) and \( q\bar{q} \rightarrow \tilde{g}\tilde{g} \) are considered. However, for \( m_{\tilde{g}} \leq 10 \text{ GeV/c}^2 \), ‘indirect’ higher-order production mechanisms in which a gluon splits to yield a \( \tilde{g}\tilde{g} \) pair becomes the dominant contribution, shown as the dashed line in Fig. 3.5. A technical problem is that because of the low efficiency for small \( m_{\tilde{g}} \), it is difficult to generate enough Monte Carlo statistics to compute the expected rates reliably. Cuts had to be placed on the Monte Carlo event generation (for ‘direct’ and ‘indirect’ gluino
Fig. 3.5 Predicted contribution to the $\Delta \phi < 140^\circ$ multijet + $E_T$ sample from supersymmetric particle production as a function of gluino mass, calculated for the case where the squark mass becomes infinitely large. The dashed curve shows the contribution from $p\bar{p} \rightarrow gX$, $g \rightarrow \tilde{g}\tilde{g}$ only [3.13].

production). This would therefore underestimate the true rates. Other assumptions which tended to underestimate the true rate included the disregard of any K-factor corresponding to higher-order QCD effects, and the neglect of other indirect production processes.

In the case when $m_\tilde{q} \ll m_\tilde{g}$, the acceptance for $\tilde{q}\tilde{q}$ final states is about a factor of three higher than for direct $\tilde{g}\tilde{g}$ production in the limit considered above. However, the cross-sections are smaller by a similar factor, so that the expected event rates finish up being comparable, as seen in Fig. 3.6. Since $e^+e^-$ data exclude $m_\tilde{q} \leq 21.5$ GeV/c$^2$, only the region of Fig. 3.6 where $m_\tilde{g} \approx 20$ GeV/c$^2$ was considered.

Fig. 3.6 Predicted contributions to the $\Delta \phi < 140^\circ$ multijet + $E_T$ sample from supersymmetric particle production as a function of squark mass, calculated for the case when the gluino mass becomes infinitely large [3.13].
Fig. 3.7 Limits on squark and gluino masses (90% CL) obtained by UA1. The arrows indicate the asymptotic values of the 90% CL contours as the squark or gluino mass becomes infinitely large [3.13].

The domain of the \((m_{\tilde{g}}, m_{\tilde{q}})\) plane excluded by this UA1 analysis is shown in Fig. 3.7. Particular cases of interest are

\[
\begin{align*}
    m_{\tilde{q}} &> 45 \text{ GeV/c}^2 \text{ in the limit } m_{\tilde{g}} \to \infty \quad (90\% \text{ CL}), \quad (3.10a) \\
    m_{\tilde{g}} &> 53 \text{ GeV/c}^2 \text{ in the limit } m_{\tilde{q}} \to \infty \quad (90\% \text{ CL}), \quad (3.10b) \\
    \tilde{h} &> 75 \text{ GeV/c}^2 \text{ if } m_{\tilde{q}} = m_{\tilde{g}} = \tilde{h} \quad (90\% \text{ CL}). \quad (3.10c)
\end{align*}
\]

Note that the detailed shape of the exclusion contour close to the diagonal is difficult to determine because the branching ratios of the \(\tilde{g}\) and \(\tilde{q}\) decays vary rapidly. The effect of a non-zero photino mass on the derived limits was studied, resulting in an insensitivity to photino masses up to \(\sim 20 \text{ GeV/c}^2\). For \(m_{\tilde{g}} > 20 \text{ GeV/c}^2\), however, the squark and gluino mass limits decrease rapidly.

The overall picture of Fig. 3.7 needs to be completed by a more detailed analysis of the light gluino case. The dominant mechanism for gluino decay is via virtual squarks, giving \(t_{\tilde{g}} \approx 4 \times 10^{-8} \text{ s } \times (m_{\tilde{g}}/m_{\tilde{q}})\), with \(m_{\tilde{q}}\) in TeV/c\(^2\) and \(m_{\tilde{g}}\) in GeV/c\(^2\) [3.14]. This means that the gluino becomes long-lived when \(m_{\tilde{g}}\) is large (e.g. \(t_{\tilde{g}} \gg 10^{-10} \text{ s}\) for \(m_{\tilde{g}} = 4 \text{ GeV/c}^2\) if \(m_{\tilde{g}} \gg 1 \text{ TeV/c}^2\)). When \(t_{\tilde{g}} \gg 10^{-10} \text{ s}\), a significant fraction of the produced gluinos decay before reaching the calorimeter, so the \(E_\text{T}\) signature is lost. However, other searches for long-lived hadrons can exclude gluinos with \(t_{\tilde{g}} \gg 10^{-10} \text{ s}\), under certain assumptions about the gluino-hadron...
interaction cross-section [3.12]. A compilation of limits for small $m_{\tilde{g}}$ and large $m_{\tilde{q}}$ is shown in Fig. 3.8. Note that there are two regions not yet excluded: i) domain I, where $m_{\tilde{g}} < 0.7$ GeV/c$^2$, and ii) domain II, where $m_{\tilde{g}} = (2.7$ to $4)$ GeV/c$^2$ and $m_{\tilde{q}} = (70$ to $800)$ GeV/c$^2$. Domain III (dotted area) is only excluded for certain assumptions on the gluino interaction cross-section [3.12]. Further dedicated searches to exclude these remaining domains are desirable, but the p$\bar{p}$ Collider is not the best machine for this task.

The UA2 Collaboration has also searched for squark and gluino production followed by $\tilde{q} \rightarrow q\tilde{\chi}$ and $\tilde{g} \rightarrow q\tilde{q}\tilde{\chi}$ decays [3.15]. Two cases were considered: i) the photino is the LSP and is therefore stable and ii) the photino decays into $\gamma\tilde{\chi}$. We first discuss the analysis performed under the assumption that the photino is stable. In this case, as mentioned above, the expected supersymmetric signature would be multijets with significant $E_T$. The background from QCD multijet events is expected to be large, particularly since large $E_T$ can be faked in UA2 by a jet escaping at polar angles $\theta$ smaller than $20^\circ$ with respect to the beams, or depositing only a small fraction of its energy in the forward calorimeters, which have only 1.5 absorption lengths. Therefore, rather stringent cuts had to be applied to remove this background. Two- or three-jet events were selected with $E_T > 12$ GeV, in which at least one pair of jets had $\Delta\phi(j_1, j_2) < 95^\circ$ to suppress the QCD background which peaks at $\Delta\phi = 180^\circ$, and $\Delta\phi(j_1, j_2) > 15^\circ$ to suppress beam-halo events. Additional technical cuts had to be applied for jets entering the forward calorimeter: the total energy measured in the forward
calorimeter had to be less than 12 GeV, and $\Delta \phi$ (total $p_T$ of central jets, $\Sigma E_T$ of forward calorimeter cells) $< 135^\circ$.

A total of 203 two- or three-jet events were selected with $E_T > 30$ GeV. To study possible $\bar{q}$ or $\bar{g}$ production, only the subsample of 13 three-jet events with $E_T > 30$ GeV were retained. An additional cut on the jet transverse energy of the third jet was applied ($E_T^j < 40$ GeV), tailored to the expected event configuration of supersymmetric processes. One event remained after these selection cuts, which could be interpreted as a $(W \rightarrow e\nu) + 2$ jets event where one of the three jets is identified as an electron. Thus UA2 quote an upper limit of

$$\sigma < 8 \text{ pb} \quad (90\% \text{ CL})$$

(3.11)
on events meeting all the above criteria. This upper limit can be used to exclude the domains of ($m_{\tilde{t}}$, $m_{\tilde{q}}$) shown in Fig. 3.9a.

The second search for squarks and gluinos was done under the assumption that the photino decays into $\gamma\tilde{H}$ (3.16). In this case, one expects the final state to contain two photons, two to six jets (depending on the production subprocess and the $\tilde{g}$, $\tilde{q}$ masses), and small $E_T$ (from the two $\tilde{H}$ escaping detection). A total of 309 events were selected with at least two photons of $E_T > 6$ GeV and $m_{\gamma\gamma} > 10$ GeV/c$^2$, most of which are single or multiple

![Fig. 3.9 Limits on squark and gluino masses (90% CL) obtained by UA2 (3.15): a) in the case of a stable photino, using three-jet events with large $p_T$; b) in the case of an unstable photino, using photon pairs accompanied by jets.](image_url)
π⁰ events. The requirement of \( m_{\gamma\gamma} > 10 \text{ GeV/c}^2 \) and at least two additional jets with \( E_T > 10 \text{ GeV} \) left UA2 with no events, corresponding to an observed cross-section of

\[
\sigma < 12 \text{ pb (90\% CL).} \quad (3.12)
\]

From a Monte Carlo simulation using the cross-sections from Ref. [3.14], UA2 inferred an overall efficiency of \( 10^{-3} \) to \( 10^{-1} \) for \( m_\tilde{g}, m_\tilde{q} \) in the range 10 to 60 GeV/c\(^2\). Combined with the limit (3.12), this led to the exclusion of the domain shown in Fig. 3.9b. Particular cases of interest are

\[
m_\tilde{g} < 9 \text{ GeV/c}^2 \text{ or } > 46 \text{ GeV/c}^2 \text{ in the limit } m_\tilde{g} \to \infty, \quad (3.13a)
\]

\[
m_\tilde{g} < 15 \text{ GeV/c}^2 \text{ or } > 50 \text{ GeV/c}^2 \text{ in the limit } m_\tilde{g} \to \infty, \quad (3.13b)
\]

\[
m_\tilde{q} \geq 60 \text{ GeV/c}^2 \text{ if } m_\tilde{g} = m_\tilde{q} = \bar{m}. \quad (3.13c)
\]

In the region of small \( \tilde{g} \) and \( \tilde{q} \) masses an uncertainty of a few GeV/c\(^2\) on the limits is expected from fragmentation effects and higher-order terms. The limits in the high-mass region are insensitive to \( m_\tilde{g} \leq 30 \text{ GeV/c}^2\).

The quoted results do not depend on the exact value of the photino lifetime, provided \( m_\tilde{\eta} > 1 \text{ GeV/c}^2\), corresponding to an average decay path in UA2 smaller than \( \sim 1 \text{ cm} \).

3.2 Sleptons

Unsuccessful searches for \( e^+e^- \to \tilde{\ell}^+\tilde{\ell}^- \), \( \tilde{\mu}^+\tilde{\mu}^- \), and \( \tilde{\tau}^+\tilde{\tau}^- \) establish the lower limits [3.17]

\[
m_\tilde{\ell} > 22 \text{ GeV/c}^2, \quad m_\tilde{\mu} > 20 \text{ GeV/c}^2, \quad m_\tilde{\tau} > 17 \text{ GeV/c}^2. \quad (3.14)
\]

The previously mentioned searches for \( e^+e^- \to \gamma + \text{nothing} \) also impose significant constraints in the \((m_\tilde{\ell}, m_\tilde{\eta})\) plane, implying, for example, that

\[
m_\tilde{\ell} > 67 \text{ GeV/c}^2 \quad (90\% \text{ CL}) \quad (3.15)
\]

if \( m_\tilde{\eta} = 0 \). However, they give no bound at all on \( m_\tilde{\ell} \) if \( m_\tilde{\eta} > 13 \text{ GeV/c}^2 \) [3.17].

The most important slepton production mechanisms at the CERN p\(\bar{p}\) Collider are \( W^* \to \tilde{\ell}^*\tilde{\nu} \) decay [3.18] and \( Z \to \tilde{\ell}^+\tilde{\ell}^- \) decay [3.19]. In each case, \( \tilde{\nu} \to \ell^*\gamma \) is expected to dominate, whilst even if the \( \tilde{\nu} \) is not the LSP and hence unstable, in many models its dominant decay is the invisible \( \tilde{\nu} \to \nu\gamma \). Therefore we expect \( \ell^* + \tilde{\eta} \) final states from \( W^* \to \tilde{\ell}^*\tilde{\nu} \) decay, and \( \ell^+\ell^- + \tilde{\eta} \) final states from \( Z \to \tilde{\ell}^+\tilde{\ell}^- \) decay. The former may be distinguishable from
conventional $W^* \rightarrow l^* \nu$ decay by having a softer $l^*$ spectrum with a different angular distribution. The $Z \rightarrow \bar{l}l$ decay can be distinguished from the conventional Drell-Yan background by its $E_T$ signature.

The UA1 Collaboration has searched for $p\bar{p} \rightarrow W + X$ followed by $W \rightarrow \bar{e}e, \bar{\nu} \rightarrow e\bar{\nu}$ decays [3.20]. It was assumed in the analysis that the $\nu$ is stable and $m_{\nu} \ll m_{W}$. Left- and right-handed sleptons are mass-degenerate and the $\bar{W}$ and $\bar{Z}$ are sufficiently heavy to be ignored. The signature sought was an excess of events with a transverse mass of $e$ and $E_T$ lower than those for the standard $W \rightarrow e\nu$ decay, and with the isotropic angular distribution in the $W$ centre of mass characteristic of spin-0 particles, to be distinguished from the sharply forward-peaked distribution characteristic of fermions with $V \pm A$ couplings.

The total $W \rightarrow e\nu$ data sample was used in the search for supersymmetric decays of the $W$, including the following requirements: $E_T > 15 \text{ GeV}$, $E_T > 15 \text{ GeV}$, and validation cuts on electron signature and $E_T$. The event sample was compared to the expectations from known processes using the ISAJET Monte Carlo predictions including full simulation of the detector. The percentages of events expected from background sources are: $W \rightarrow e\nu$ (92%), $W \rightarrow \tau\nu$ (6%), and QCD jet fluctuations (2%). The transverse mass and angular distribution observed are well described by the above-mentioned background contributions. Using this agreement

![Figure 3.10](image)

and the expected distributions from a supersymmetric decay of the $W$, a likelihood fit was performed in the $m_T - q \cos \theta$ plane to extract the excluded domain shown in Fig. 3.10. Of particular interest is the bound *)

$$m_\xi > 32 \text{ GeV}/c^2 \quad (90\% \text{ CL})$$  \hspace{1cm} (3.16)

*) This limit represents an improvement on the previous result [3.20] and is due to the increased statistics from the 1985 run and to a more detailed Monte Carlo study of the supersymmetric processes.
in the case that \( m_{\tilde{e}} = m_{\tilde{\mu}} \). Note that this limit applies only to the \( \tilde{e}_L \), since the W does not couple to the \( \tilde{e}_R \).

The UA1 Collaboration has also searched for the decay of the \( W \to \tilde{\ell}\tilde{\ell} \) and the \( Z \to \tilde{\ell}\tilde{\ell} \) (\( \tilde{\ell} = \tilde{e} \) or \( \tilde{\mu} \)) with \( \tilde{\ell} \to \ell \gamma \). The \( Z \to \tilde{e}\tilde{e} \) mode is less sensitive than is the \( W \to \tilde{\ell}\tilde{\ell} \) channel described above, resulting in a weaker limit on \( m_{\tilde{e}} \). In the present data of the muonic decay channels, the sensitivity is negligible in the \( Z \) case and marginal in the \( W \) case owing to the inferior momentum resolution (compared with the energy resolution), resulting in a broader transverse or invariant mass distribution.

The UA2 Collaboration has searched for \( p\bar{p} \to Z + X \) followed by \( Z \to \tilde{e}^+\tilde{e}^- \), \( \tilde{e}^+ \to e^+\gamma \) decays [3.15]. In this analysis it was assumed that the \( \tilde{\gamma} \) is stable (therefore escaping detection), that \( m_{\tilde{e}_L} = m_{\tilde{e}_R} \), and that \( m_{\tilde{e}} > m_Z/2 \). This supersymmetric process is characterized by final states containing low-mass electron pairs and missing transverse energy \( E_T \). The total data sample from the \( Z \) triggers, corresponding to an integrated luminosity of 910 nb\(^{-1}\), has been used in this search. The standard UA2 \( Z \) analysis has been modified in such a way that stringent electron validation cuts have been applied to both electron candidates in order to reduce the dominant background from fake electron pairs at low values of the electron-pair mass. Figure 3.11 shows a scatter plot of \( m_{e^+e^-} \) versus \( p_T^Z \) for

![Figure 3.11](image)

Fig. 3.11 Distribution of the electron-pair mass versus \( p_T^Z \) for the 57 selected events. This sample contains 21 \( Z \to e^+e^- \) candidates and 16 \( \pm 6 \) Drell–Yan pairs after background subtraction. Also shown is the dashed area where most of the events coming from \( Z \to \tilde{e}\tilde{e} \) are expected [3.15].
the total of 57 electron-pair candidates selected with \( m_{ee} > 10 \text{ GeV/c}^2 \). The variable \( p_T^e \) is the projection of \( p_T^{e^-} \) on the bisector of the azimuthal angle between the two electron transverse momenta. Note that although \( \ell_T \) is in principle a signature of this process, it could not be used in the UA2 analysis because of the incomplete angular coverage of the calorimeter. The data with \( m_{e^+e^-} < 70 \text{ GeV/c}^2 \) comprized 36 events, of which \( 20 \pm 1 \) were estimated to be QCD background. The signal from genuine electron pairs is therefore \( 16 \pm 6 \) events for \( m_{e^+e^-} < 70 \text{ GeV/c}^2 \). A Monte Carlo simulation was used to estimate the contribution from \( Z \rightarrow \ell\bar{\ell} \) decay to the selected data sample. A likelihood fit to the \((m_{e^+e^-}, p_T^e)\) distribution of Fig. 3.11 was performed, taking into account the expected contributions from the Drell–Yan continuum, from QCD background, and from \( Z \rightarrow \ell\bar{\ell} \) decays. This analysis resulted in the 90% CL limits in the \((m_\ell, m_\gamma)\) plane shown in Fig. 3.12 together with limits from \( e^+e^- \) experiments [3.17]. The shaded area in Fig. 3.12 indicates the variation of the limit if the expected Drell–Yan contribution is increased by 50%.

Fig. 3.12 Limits (90% CL) on the \( \bar{\ell} \) and \( \gamma \) masses obtained by UA2 [3.15], assuming a stable photino. The dashed area indicates the theoretical uncertainty in Drell–Yan production arising from higher-order terms. Also shown are recent limits from \( e^+e^- \) experiments.

3.3 Winos and Zinos

Unsuccessful searches for \( e^+e^- \rightarrow \tilde{\nu}\tilde{\nu} \) tell us that [3.17]

\[
m_{\tilde{\nu}^\pm} \gg 22 \text{ GeV/c}^2.
\]

whilst unsuccessful searches for \( e^+e^- \rightarrow \gamma Z \) give us bounds on \( m_Z \) which range up to 40 GeV/c² [3.17], but depend on the values assumed for \( m_\gamma \) and \( m_\tilde{\nu} \). If \( m_{\tilde{\nu}^\pm} \ll m_{\tilde{\nu}^0} \), we expect decay modes and branching ratios similar to those for conventional charged and neutral heavy leptons. If \( m_\gamma < m_{\tilde{\nu}^\pm} < m_{\tilde{\nu}^0} \), we expect the decays \( \tilde{\nu}^\pm \rightarrow \ell^\pm \nu \) and \( \ell^\pm \bar{\nu} \) and
\( \bar{Z} \rightarrow \ell^+ \ell^* \) and \( \tilde{\nu} \) to dominate. If \( m_{\tilde{g}} < m_{\tilde{W}^*} \), we expect the decays \( \bar{W}^* \rightarrow \tilde{\nu} \ell^* \), \( \tilde{q} \bar{q} \), etc., and \( \bar{Z} \rightarrow \ell^\nu \tilde{\nu} \), \( \tilde{q} \bar{q} \), etc., to dominate, with branching ratios similar to those of the \( W^* \) and \( Z \) into the corresponding particles. The relative magnitudes of \( m_{\tilde{\nu}} \), \( m_{\tilde{W}^*} \), and \( m_{\tilde{q}} \) are uncertain, although in most models \( m_{\tilde{\nu}} < m_{\tilde{q}} \).

The dominant production mechanisms for the \( \bar{W}^* \) and \( \bar{Z} \) are likely to be:

- \( \bar{W}^* \rightarrow \bar{W}^* \tilde{\tau} \) giving \( \ell^* + \ell^\tau \), or dijet + \( \ell_T \) events;
- \( \bar{W}^* \rightarrow \bar{W}^* \tilde{\tau} \) giving \( 3 \ell^* + \ell^\tau \), or dijet + \( \ell_T \), or \( \ell^* + \ell^\tau \) + dijet + \( \ell_T \) events;
- \( \tilde{\tau}^* \tilde{\tau}^- \) + dijet + \( \ell_T \), or quadrijet + \( \ell_T \) events;
- \( \bar{Z} \rightarrow \bar{W}^+ \bar{W}^- \) giving \( \ell^* \tilde{\tau}^- + \ell^\tau \), or \( \ell^* + \ell^\tau \), or quadrijet + \( \ell_T \) events;
- \( \bar{Z} \rightarrow \bar{Z} \bar{Z} \) giving \( 4 \ell^* + \ell_T \), or \( 2 \ell^* + \ell_T \), or quadrijet + \( \ell_T \) events.

Clearly, not all these jets and/or leptons would be distinguished by experiment, so we must look for \( \bar{W}^* \) and \( \bar{Z} \) in a combination of \( n \ell^* + m \) jets + \( \ell_T \) event classes with \( n \leq 4 \) and \( m \leq 4 \).

The only published limit from this class of searches is by UA2 [3.15], looking for \( e^+e^- \) events with non-zero \( p_T^\ell \), as for the slepton search described in subsection 3.2. A Monte Carlo simulation was performed to estimate the expected contribution from \( Z \rightarrow \bar{W}W \rightarrow e^+e^- \bar{\nu}\bar{\nu} \) decays. It was assumed that the branching ratio for \( \bar{W} \rightarrow e\bar{\nu} \) is 33%, that the \( \tilde{\nu} \) is invisible, and that \( m_{\tilde{\nu}} > m_{\tilde{W}^*}/2 \). The dashed area in Fig. 3.11 indicates where most of the events coming from \( Z \rightarrow \bar{W}W \) are expected. A likelihood fit to the distribution in the \( (m_{e^+e^-}, p_T^\ell) \) plane analogous to that described in subsection 3.2 led to the excluded domain of the \( (m_{W^*}, m_{\tilde{s}}) \) plane shown in Fig. 3.13 together with limits from \( e^+e^- \) experiments [3.17]. The shaded area in

![Graph](image_url)

Fig. 3.13 Limits (90% CL) on the \( \bar{W}^* \) and \( \tilde{s} \) masses obtained by UA2 [3.15]. The dashed area indicates the theoretical uncertainty in Drell–Yan production arising from higher-order terms. Also shown are recent limits from \( e^+e^- \) experiments.
Fig. 3.13 shows the change in the limit if the expected Drell–Yan contribution is increased by 50%, taking into account the larger higher-order QCD corrections. In the W̅ case, the UA2 results improve significantly the values obtained in e⁺e⁻ experiments so far.

4. ADDITIONAL GAUGE BOSONS

There are many proposals for additional gauge bosons. Here we mention just a few so as to suggest the range of possibilities. The SU(2)ₖ × U(1)ᵧ gauge group of the Standard Model is not particularly handsome, nor has it been derived convincingly from any TOE. Therefore it is natural to consider expansions of the gauge group which might i) make it more elegant, or ii) relate it more directly to some TOE framework. Alternatively, some authors hope to solve the problem of the origins of the W⁺ and Z mass in a composite model, in which case other more massive vector bosons can be expected.

Typical of the first approach are proposals that the gauge group be expanded to SU(2)ₖ × SU(2)₁🔗 × U(1) with parity violated spontaneously [4.1]. In this case one would expect three additional gauge bosons—two charged bosons W⁺ᵧ and one neutral boson Z'. In this framework it is natural to expect that the magnitudes of the W⁺ᵧ couplings could be similar to those of the known W. Also, in this approach the Z' would have couplings similar to the Standard Model Z. There are important bounds on the mass of the W⁺ᵧ from μ decay [4.2],

\[ W⁺ᵧ > 432 \text{ GeV/c}² \quad (90\% \text{ CL}), \]  \( (4.1) \)

and from the magnitude of K⁰–K̅⁰ mixing [4.3],

\[ W⁺ᵧ > 1.6 \text{ TeV/c}² \quad (90\% \text{ CL}). \]  \( (4.2) \)

However, the first of these bounds assumes the simultaneous presence of a light right-handed neutrino, which may not be true. The second bound makes assumptions regarding the nature of generalized right-handed Cabibbo mixing which are plausible but not ironclad.

Typical of the TOE approach to augmentation of the gauge group are four-dimensional models inspired by the superstring [4.4]. Early efforts in this direction compactified the ten-dimensional heterotic string with an E₅ × E₅ gauge group on a Calabi–Yau manifold to obtain a four-dimensional gauge group which was some subgroup of E₅. The minimal possibility was SU(2)ₖ × U(1)ᵧ × U(1)ₑ; larger possibilities such as SU(2)ₖ × U(1)ᵧ × U(1)¹ or SU(2)ₖ × SU(2)ₑ × U(1)¹ or² also exist. It is a generic feature of such models [4.5] that the U(1) gauge couplings are expected to have a magnitude similar to that of the U(1)ᵧ in the Standard Model, so that gᵧ/gₑ = tan²θₑ, whilst the SU(2)ₑ gauge couplings are of order gₑ as usual. The extra U(1) hypercharges Yₑ of the known particles in the minimal SU(2)ₖ × U(1)ᵧ × U(1)ₑ extension of the Standard Model (which we call Model A) are completely fixed.
Table 4.1
Possible neutral currents in superstring models

<table>
<thead>
<tr>
<th></th>
<th>$T_{3L}$</th>
<th>$\sqrt{(b/3)}Y'$</th>
<th>$\sqrt{(b/3)}Y' - \sqrt{(b/3)}Y''$</th>
<th>$\sqrt{(b/3)}Y' + \sqrt{(b/3)}Y''$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$(u, d)_L$</td>
<td>$\pm \frac{1}{2}$</td>
<td>$\frac{1}{6}$</td>
<td>$\frac{1}{3}$</td>
<td>0</td>
</tr>
<tr>
<td>$u_L$</td>
<td>0</td>
<td>$-\frac{2}{3}$</td>
<td>$\frac{1}{3}$</td>
<td>0</td>
</tr>
<tr>
<td>$d_L$</td>
<td>0</td>
<td>$\frac{1}{3}$</td>
<td>$-\frac{1}{6}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$(\nu, \ell)_L$</td>
<td>$\pm \frac{1}{2}$</td>
<td>$-\frac{1}{2}$</td>
<td>$-\frac{1}{6}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$e_L$</td>
<td>0</td>
<td>1</td>
<td>$\frac{1}{3}$</td>
<td>0</td>
</tr>
<tr>
<td>$D_L$</td>
<td>0</td>
<td>$-\frac{1}{3}$</td>
<td>$-\frac{2}{3}$</td>
<td>0</td>
</tr>
<tr>
<td>$D^\ell_L$</td>
<td>0</td>
<td>$\frac{1}{3}$</td>
<td>$-\frac{1}{6}$</td>
<td>$-\frac{1}{2}$</td>
</tr>
<tr>
<td>$\nu_L$</td>
<td>0</td>
<td>0</td>
<td>$\frac{5}{6}$</td>
<td>$-\frac{1}{2}$</td>
</tr>
<tr>
<td>$N_L$</td>
<td>0</td>
<td>0</td>
<td>$\frac{5}{6}$</td>
<td>$\frac{1}{2}$</td>
</tr>
<tr>
<td>$(H^+, H^0)_L$</td>
<td>$\pm \frac{1}{2}$</td>
<td>$\frac{1}{2}$</td>
<td>$-\frac{1}{2}$</td>
<td>0</td>
</tr>
<tr>
<td>$(\tilde{H}^+, \tilde{H}^0)_L$</td>
<td>$\pm \frac{1}{2}$</td>
<td>$-\frac{1}{2}$</td>
<td>$-\frac{1}{6}$</td>
<td>$-\frac{1}{2}$</td>
</tr>
</tbody>
</table>

Model A: $Z'$ couples to $Y'$
Model B: $Z'$ couples to $[\sqrt{(b/3)}Y' - \sqrt{(b/3)}Y'']$
Model C: $Z'$ couples to $[\sqrt{(b/3)}Y' + \sqrt{(b/3)}Y'']$

and are shown in the first column in Table 4.1. In models based on $SU(2)_L \times U(1)_Y \times U(1)^2$ it is often postulated that one of the additional $U(1)$ factors is spontaneously broken by a large field $v_{e.v.}$ and is hence irrelevant to contemporary phenomenology. The remaining orthogonal $U(1)$ combination is model-dependent, and two possibilities have been favoured in the literature, which we denote as Models B and C [4.6]. The extra $U(1)$ hypercharges in these two cases are listed in the second and third columns of Table 4.1. Bounds on the masses of these various $U(1)$ bosons can be obtained from low-energy $\nu\nu$, $\nu e$, and $e e$ scattering, from $e^+e^-$ annihilation, and from the observed values of the $W^*$ and $Z$ masses. Making a global fit to all these electroweak data but without imposing all the constraints on the vector boson masses which might be expected from the conventional Higgs structure, it has been found that [2.22]

\[
m_{Z'} > \left\{ \begin{array}{c}
\text{Model A:} & 129 \text{ GeV/c}^2 \\
\text{Model B:} & 352 \text{ GeV/c}^2 \\
\text{Model C:} & 180 \text{ GeV/c}^2 \\
\end{array}\right\} \quad 90\%\, \text{CL.}
\]

We will briefly mention how these limits compare with those from the CERN p$\bar{p}$ Collider. An important ambiguity in obtaining these comes from the ignorance of the $Z'$ decay modes. The $Z'$ may be able to decay into exotic particles and sparticles which are not kinematically accessible to conventional $W^*$ and $Z$ decays. This could suppress the observable $Z' \rightarrow e^+e^-$.
Table 4.2

$Z'$ width and branching ratio into $e^+e^-$ pairs for three representative 'superstring-inspired' models: (i) $Z'$ can decay into only the observed fermions and the top quark; (ii) $Z'$ can decay into three families of 27 fermions and their supersymmetric partners [4.6]

<table>
<thead>
<tr>
<th>Model</th>
<th>Case</th>
<th>$10^3 \times \Gamma_{Z'}/m_{Z'}$</th>
<th>BR$(e^+e^-)$%</th>
</tr>
</thead>
<tbody>
<tr>
<td>A</td>
<td>I</td>
<td>6.5</td>
<td>3.6</td>
</tr>
<tr>
<td></td>
<td>II</td>
<td>38</td>
<td>0.6</td>
</tr>
<tr>
<td>B</td>
<td>I</td>
<td>12</td>
<td>5.9</td>
</tr>
<tr>
<td></td>
<td>II</td>
<td>38</td>
<td>1.8</td>
</tr>
<tr>
<td>C</td>
<td>I</td>
<td>6.5</td>
<td>5.4</td>
</tr>
<tr>
<td></td>
<td>II</td>
<td>38</td>
<td>0.9</td>
</tr>
</tbody>
</table>

or $\mu^+\mu^-$ branching ratio by a significant factor, as shown in Table 4.2. Because of the experimental limits on the total $W^*$ and $Z$ decay widths reported elsewhere in this volume [2.3], there are constraints on the masses of such exotic particles which could in principle give interesting lower bounds on the $Z' \rightarrow e^+e^-, \mu^+\mu^-$ branching ratios. In practice, however, branching ratios almost as low as the minima shown in Table 4.2 are possible. When quoting lower limits on superstring $Z'$ masses, we will consider both the minimal and maximal cases in Table 4.2.

Both UA1 and UA2 have published negative searches for additional gauge bosons $W'$ and/or $Z'$. In the case of UA1 [4.7], no $W' \rightarrow e\bar{\nu}_e$ candidates have been observed with $m_{W'}$ beyond the distribution expected for conventional $W \rightarrow e\bar{\nu}_e$ decays, corresponding to

$$\alpha \cdot \text{BR}(W' \rightarrow e\bar{\nu}_e) < 4.6 \text{ pb} \quad (90\% \text{ CL}). \quad (4.4)$$

If we assume that the $W'$ has couplings to quarks of the same magnitude as the conventional $W$ and that the $W' \rightarrow e\nu_e$ branching ratio is the same as the $W \rightarrow e\nu_e$ branching ratio, and use the quark distributions of Diemoz et al. [4.8], to estimate the $W'$ production cross-section, then from formula (4.4) we can deduce that [4.7]

$$m_{W'} > 220 \text{ GeV}/c^2. \quad (4.5)$$

This limit should be directly applicable to the $W_R$ of an SU(2)$_L \times$ SU(2)$_R \times U(1)_{Y}$ model. In the case of $Z' \rightarrow l^+l^-$, no $e^+e^-$ pair was found beyond the distribution expected for $Z \rightarrow e^+e^-$ decays. On the basis of these negative results, it was concluded that [4.7]
\[ \sigma \cdot \text{BR}(Z' \to e^+e^-) < 4.7 \text{ pb} \quad (90\% \text{ CL}). \quad (4.6) \]

Using the structure functions of Diemoz et al. \cite{4.8} and assuming that the $Z'$ has quark couplings of similar strength to those of the $Z$, as well as a similar $e^+e^-$ branching ratio, it was concluded that \cite{4.7}

\[ m_{Z'} > 173 \text{ GeV/c}^2 \quad (90\% \text{ CL}). \quad (4.7) \]

This limit could be directly applicable to a neutral gauge boson in an SU(2)$_L \times$ SU(2)$_R \times$ U(1)$_Y$ model, but is not applicable to superstring-inspired models which have smaller couplings and possibly smaller leptonic branching ratios.

In Fig. 4.1 the cross-section times branching ratio ($\sigma$-BR) is plotted as a function of the additional vector boson masses. The solid line (dashed line) corresponds to the limit obtained using structure functions of Diemoz et al. (Duke and Owens set 1), where the curves are

![Graph showing limits on additional vector boson masses](image)

**Fig. 4.1** Limits (90% CL) obtained on ($\sigma$-BR) as a function of additional vector-boson masses at $\sqrt{s} = 630$ GeV \cite{4.7} using structure functions \cite{4.8} of Diemoz et al. (solid line) and of Duke and Owens (dashed line). The curves are normalized to the UA1 results at the $W$ and $Z$ mass.
normalized to the UA1 results at the W and Z mass. The muonic decay channels have not been used in the searches for additional vector bosons because of the inferior momentum resolution of the measured muon track; the resulting transverse and invariant mass distributions would give a much weaker limit.

The UA2 Collaboration has performed similar analyses [3.15], but has quoted directly lower bounds on $m_{W'}$ and $m_{Z'}$ as functions of $\lambda^2$, which is defined by

$$\lambda^2 = \sum_q g_{q\bar{q}Z'}^2/\sum_q g_{q\bar{q}Z}^2, \quad B_e = BR(Z' \rightarrow e^+e^-)/BR(Z \rightarrow e^+e^-).$$ (4.8)

Four different data samples have been used to search for additional vector bosons:

i) for $m_{W'} > m_W$, the standard UA2 sample of $W \rightarrow e\nu$ candidates with $p_T > 20$ GeV/c and $m_T(e\nu) > 50$ GeV/c$^2$;

ii) for $m_{W'} < m_W$, electron candidates with $p_T > 12$ GeV/c and no jet activity opposite in azimuth to the electron; the estimated contribution from the dominant background coming from two-jet events, where one jet fakes the electron signature and the other jet escapes detection, is included in the final limit;

iii) for $m_{Z'} > m_Z$, the standard UA2 sample of $Z \rightarrow e^+e^-$ candidates with $m_{ee} > 76$ GeV/c$^2$;

iv) for $m_{Z'} < m_Z$, the low-mass electron-pair sample as described in subsection 3.2.

A Monte Carlo simulation was performed for $W'$ and $Z'$ production and decay using the structure functions of Duke and Owens. Likelihood fits to the data samples were then performed, using the electron transverse momentum spectrum for the case of $W'$ and the invariant mass spectrum of the two electrons for the case of $Z'$.

The UA2 90% CL exclusion contours for $W' \rightarrow e\bar{\nu}_e$ and $Z' \rightarrow e^+e^-$ are shown in Figs. 4.2a and 4.2b. This search is insensitive to a $W'$ and/or $Z'$ almost degenerate in mass with the Standard Model vector bosons, because large uncertainties in theoretical predictions and experimental cross-section measurements prevent the exclusion of an additional contribution above the expected rates for the $W$ and the $Z$. If we assume that the $W'$ has the same couplings and $BR(W' \rightarrow e\bar{\nu}_e)$ as the Standard Model, we can read off from Fig. 4.2a that [3.15]

$$m_{W'} > 209 \text{ GeV/c}^2 \quad (90\% \text{ CL}).$$ (4.9)

The same assumption for the $Z'$ leads to [3.15]

$$m_{Z'} > 180 \text{ GeV/c}^2 \quad (90\% \text{ CL}),$$ (4.10)

as can be seen from Fig. 4.2b. An additional $W'$ with $m_{W'} < 25$ GeV/c$^2$ and an additional $Z'$ with $m_{Z'} < 50$ GeV/c$^2$ would have been detected at the ISR and/or PETRA. Therefore additional vector bosons with masses smaller than the standard ones are excluded for
Fig. 4.2 Limits (90% CL) on additional vector bosons W' and Z' depending on $\lambda_5^2 B_e$ (coupling times branching ratio) obtained by UA2 [3.15]: a) for the W' obtained from an analysis of single electron candidates; b) for the Z' obtained from an analysis of electron pairs.

couplings to quarks and leptons similar to the Standard Model ones. For future reference we note that the exclusion contour in Fig. 4.2b and the special case (4.10) correspond, in general, to

$$\sigma \cdot \text{BR}(Z' \rightarrow e^+e^-) < 4 \text{ pb} \quad (90\% \text{ CL}). \quad (4.11)$$

As an example of a model with $\lambda_5^2 B_e$ (4.8) less than unity, let us consider what may occur in the minimal rank-5 superstring-inspired model, where the extra U(1) hypercharges $Y_E$ are those specified in column 1 of Table 4.1, and the coupling strength $\alpha_E = 0.016$ is specified by the renormalization group. These predictions imply that $\lambda_5^2 < 1$. As mentioned above, two alternative hypotheses can be formulated about the Z' decays: either it decays only into conventional quarks and leptons—the 'minimal' scenario; or it can decay into three generations of particles in the full 27 representations of $E_6$, together with their supersymmetric partners—the 'maximal' scenario. In the 'maximal' case, we must nevertheless respect experimental bounds on the masses of sleptons [2.13, 2.21] and other charged particles in $e^+e^-$ annihilation, as well as bounds on the squark masses from the CERN pp Collider itself [3.13, 3.15].

These particles could be light enough to appear in Z decays as well as in Z' decays. The most conservative lower bound on the Z' mass is obtained by adjusting the undiscovered
particle masses consistently with the experimental bounds, so as to fit the observed \( \sigma \cdot \text{BR}(Z \to e^+e^-) \) satisfactorily whilst minimizing \( \sigma \cdot \text{BR}(Z' \to e^+e^-) \). This has been done for the minimal rank-5 and other superstring-inspired models. Comparing the resulting minimized \( \sigma \cdot \text{BR}(Z' \to e^+e^-) \) with the upper limit obtained by combining preliminary results of UA1 and the UA2 value (4.11), we get (4.9)

\[
\sigma \cdot \text{BR}(Z' \to e^+e^-) < 1.8 \text{ pb} \quad (90\% \text{ CL}).
\]

(4.12)

The resulting lower bounds on \( m_{Z'} \) are

\[
m_{Z'} > \begin{cases} 
118 \text{ GeV}/c^2 \\
140 \text{ GeV}/c^2 \quad (90\% \text{ CL}) \\
115 \text{ GeV}/c^2 
\end{cases}
\]

(4.13)

in the ‘maximal’ decay scenario, to be compared with

\[
m_{Z'} > \begin{cases} 
167 \text{ GeV}/c^2 \\
171 \text{ GeV}/c^2 \quad (90\% \text{ CL}) \\
158 \text{ GeV}/c^2 
\end{cases}
\]

(4.14)

in the ‘minimal’ decay scenario, where the \( Z' \) decays only into conventional quarks and leptons. We see that the more conservative bounds (4.13) are less stringent than those (4.3) already available from the neutral-current data, whilst the bounds (4.14) obtained in the ‘minimal’ scenario are more stringent. The optimal conservative bound can be obtained by combining the neutral current bounds (4.3) with the ‘maximal’ decay collider bounds (4.13). This exercise has been carried out for the minimal rank-5 model, with the result (4.9)

\[
m_{Z'} > 156 \text{ GeV} \quad (90\% \text{ CL}),
\]

(4.15)

which is considerably stronger than the individual bounds for this model.

Before leaving the superstring-inspired \( Z' \), we remind the reader that because of its mixing with the Standard Model \( Z \), the \( Z' \) has in general a non-zero coupling to \( W \) pairs, which can yield

\[
\text{BR}(Z' \to W^+W^-)/\text{BR}(Z' \to \ell^+\ell^-) = O(1).
\]

(4.16)

However, as might be guessed from the present lower limits (4.13) and (4.14) on the \( Z' \) mass from searches for \( Z' \to \ell^+\ell^- \), the present data samples of UA1 and UA2 are too small to
Fig. 4.3 Total cross-section $\sigma(pp \rightarrow Z' \rightarrow W^+W^-)$ as a function of $m_{Z'}$ for Model C (see text) at $\sqrt{s} = 630$ GeV. The expected numbers of events are given for total integrated luminosities of 10 pb$^{-1}$ and 730 nb$^{-1}$ [4.6].

expect a signal for $Z' \rightarrow W^+W^-$. However, this is a possible prospect for the larger data sets obtainable with ACOL, as can be seen from the cross-sections shown in Fig. 4.3 for model C [4.6].

5. COMPOSITE MODELS

There are several theoretical motivations for composite models. One of these is the proliferation of quark and lepton flavours, which can give the idea that quarks and leptons may be composite particles [5.1]. Another is the massive nature of the $W^\pm$ and $Z$ bosons, which suggests to some theorists the possibility that they may be composite, like the $Q$ vector meson in QCD [5.2]. Alternatively, one response to the problematic quadratic divergences (3.1) in the mass of the elementary Higgs boson has been to suggest that this boson may be composite [5.3]. We are unaware of any experimental probes of this last idea at the CERN $p\bar{p}$ Collider, and we will concentrate on tests of composite models of the quarks and leptons.

There are several possible manifestations of quark and lepton compositeness, including new contact interactions having the new strong interaction scale $\Lambda$, quark and lepton form factors, and excited quarks $q^*$ or leptons $l^*$. A general framework for discussing four-fermion contact interactions has been set up by Eichten, Lane and Peskin [5.4]:

$$\mathcal{L}_{\text{ELP}} = g_{\text{eff}}^2 \left[ \eta_{LL}/2\Lambda_{LL}^2 \left( \bar{\psi}_L \gamma^\mu \psi_L \right) \left( \bar{\psi}_L \gamma^\nu \psi_L \right) + \eta_{RR}/2\Lambda_{RR}^2 \left( \bar{\psi}_R \gamma^\mu \psi_R \right) \left( \bar{\psi}_R \gamma^\nu \psi_R \right) + \eta_{LR}/2\Lambda_{LR}^2 \left( \bar{\psi}_L \gamma^\mu \psi_L \right) \left( \bar{\psi}_R \gamma^\nu \psi_R \right) + \eta_{RL}/2\Lambda_{RL}^2 \left( \bar{\psi}_R \gamma^\mu \psi_R \right) \left( \bar{\psi}_L \gamma^\nu \psi_L \right) \right],$$

(5.1)

where $\Lambda_{ij}$ is proportional to the binding scale of the composite model and we have allowed for different possible helicity structures of the four-fermion interactions ($\eta_{ij}$ may be $\pm 1$, and we assume that $g_{\text{eff}}^2 = 4\pi$). Experiments at the CERN $p\bar{p}$ Collider are only sensitive to four-quark contact interactions. Their effects may be probed either by looking at the total rate for two-jet
events, or by measuring the two-jet angular distributions: the contact terms (5.1) give much more isotropic distributions than does one-gluon exchange. The principal experimental problem in interpreting any measurement of the total two-jet cross-section is the theoretical uncertainty in the initial-state proton distributions and in higher-order QCD effects. Possible limitations on bounding \( \Lambda \) using the angular distributions are a limited angular acceptance for jets and radiative corrections to the QCD tree-level angular distributions.

There is little sensitivity to quark or lepton form factors at the CERN p\( \bar{p} \) Collider, and there has not been any experimental interest in probing them, so we will not discuss them further.

To discuss the production of excited quarks \( q^* \) and leptons \( \ell^* \) it is convenient to use the following general parametrization of the \( f^-f^+ \) gauge-boson vertex,

\[
\mathcal{L} = (g/M) \bar{f} \gamma_\mu f F^{\mu\nu},
\]

where again we expect that \( M \) is of the order of the compositeness scale. We can use formula (5.2) — where \( g \) is the SU(3) gauge coupling, \( f \) is a quark, \( F^{\mu\nu} \) is a gluon field strength, and \( f^* \) is an excited quark — to estimate the total cross-section for \( g + q \to q^* \). This cross-section can then be compared with experimental bounds on dijet mass bumps, in order to give a lower bound on \( M \) as a function of \( m_{q^*} \). We expect the dominant decay mode of the \( q^* \) to be into \( q + g \), but \( q + \gamma \), \( W^* \), and \( Z \) decays are also possible, and could have comparable rates if the corresponding \( M \) values are similar. There have also been searches for bumps in \( \gamma + \text{jet} \) combinations which could come from \( q^* \to q + \gamma \) decay, and for excited leptons in \( Z \to \ell^+\ell^- \) and \( \ell^+\ell^- \) and \( W \to \ell^+\nu^* \) and \( \ell^*\nu \) decays, stimulated by the early observations [5.5] of radiative \( Z \) decays, \( Z \to e^+e^- + \gamma \) and \( Z \to \mu^+\mu^- + \gamma \). Models for these decays using vertices of the form (5.2) have been proposed. They tend to give angular distributions that are closer to phase space than are the observed angular distributions, which are more similar to those predicted by calculation of QED radiative corrections to \( Z \to e^+e^- \) and \( Z \to \mu^+\mu^- \). There have also been attempts to interpret the radiative \( Z \) decays in the context of composite models of the vector bosons.

The experimental searches at the CERN p\( \bar{p} \) Collider for contact interactions of the form (5.1) have been moulded by the characteristics of the detectors. Because of its limited rapidity coverage, the UA2 Collaboration has concentrated on the overall normalization of the large-\( p_T \) jet cross-section [5.6]. Finite values of \( \Lambda_c \) would produce an excess of events, compared with ordinary QCD predictions (\( \Lambda_c = \infty \)) at large \( p_T \). Under the assumption that the main uncertainties (systematic errors, K-factor) are about constant over the full \( p_T \) range, deviations in the high-\( p_T \) tail can be observed. Taking into account both the theoretical and experimental uncertainties, the following limit was obtained [5.6]:

\[
\Lambda_c > 370 \text{ GeV} \quad (95\% \text{ CL}).
\]
As already commented above, this technique is limited by theoretical uncertainties in the large-$p_T$ jet cross-sections, as well as by systematic errors in the overall energy calibration of the calorimeters. To avoid the latter problem and also to take advantage of the larger rapidity coverage, UA1 preferred to use the two-jet angular distribution to bound $\Lambda_c$. Figure 5.1 shows the angular distribution observed for jet pairs with masses between 240 GeV/c$^2$ and 300 GeV/c$^2$ [5.7]. The variable $\chi$ is defined as

$$
\chi = \frac{1 + \cos \theta}{1 - \cos \theta}, \quad (5.4)
$$

where the c.m. scattering angle $\theta$ was calculated as the angle between the axis of the jet pair and the beam direction in the jet-jet rest frame [5.8]. The distribution in $\chi$ would be almost flat in the absence of perturbative QCD scaling violations. A finite parton size implying substructure will modify this angular distribution such that more events are expected at wide angles relative to the QCD prediction. The leading-order QCD calculation, including scale-breaking effects and using the EHLQ structure functions [5.9] with $\Lambda = 200$ MeV and $Q^2 = p_T^2$, fits the angular distribution very well, as can be seen from the solid line in Fig. 5.1. The

![Normalized angular distribution for very high mass jet-pairs as a function of $\chi$. The solid curve is the QCD production ($\Lambda_c = \infty$); the dotted curve corresponds to $\Lambda_c = 300$ GeV, which is clearly excluded by the UA1 data [5.7].](image)

QCD prediction corresponds to $\Lambda_c = \infty$. The dashed line is obtained with $\Lambda_c = 300$ GeV in the effective Lagrangian (5.1). By varying the value of the parameter $\Lambda_c$ in the fit to the measured angular distribution and taking into account the systematic uncertainty on the jet-energy scale, UA1 obtained the lower limit [5.7]

$$
\Lambda_c > 415 \text{ GeV} \quad (95\% \text{ CL}). \quad (5.5)
$$
This figure is less stringent than the corresponding bounds from $e^+e^-$ annihilation, which however do not probe for compositeness in the $qq\bar{q}\bar{q}$ interaction.

Some excitement was sparked among composite modellers a few years ago by the observation of three radiative $Z$ decays in the 1983 data sample: one $Z \rightarrow e^+e^-\gamma$ event from UA2, and one $Z \rightarrow e^+e^-\gamma$ and one $Z \rightarrow \mu^+\mu^-\gamma$ event from UA1 [5.5]. Key parameters of these three events are gathered in Table 5.1. Figure 5.2 shows Dalitz plot distributions for the observed three events, and also distributions expected from conventional radiative decays calculated using QED, and from some exotic sources [5.10]. The following expressions have been used for the plots in Fig. 5.2:

\[
x_L = \frac{1}{2} \left[ \text{lower m}^2(\ell\gamma)/\text{m}^2(\ell^+\ell^-\gamma) \right],
\]

\[
x_H = \frac{1}{2} \left[ \text{higher m}^2(\ell\gamma)/\text{m}^2(\ell^+\ell^-\gamma) \right],
\]

\[
x = \frac{1}{2} \left[ \text{m}^2(\ell^+\ell^-)/\text{m}^2(\ell^+\ell^-\gamma) \right],
\]

which satisfy the relation

\[
x_L + x_H + x = 1. \tag{5.7}
\]

The observed events looked qualitatively like conventional radiative decays, but the energies and/or angles of the observed photons were surprisingly high for conventional bremsstrahlung. Since the initial excitement, many more $Z \rightarrow e^+e^-$ and $Z \rightarrow \mu^+\mu^-$ events have been accumulated and only one more $Z \rightarrow \mu^+\mu^-\gamma$ candidate has been observed in the UA1 data sample. The event parameters are given in the last column of Table 5.1. Compared with the total event sample, the radiative events seen are no longer surprising from the point of view of conventional QED radiative decay calculations. The probability of observing an event like the one seen in the total UA2 event sample is now about 0.4 [5.11].

Inspired by the early observations of these radiative $Z$ decays, searches were also made for radiative $W \rightarrow \ell\nu\gamma$ events.

The UA1 Collaboration searched for massive $e\nu\gamma$ and $\mu\nu\gamma$ final states containing an energetic photon using the 1983 data sample corresponding to 136 nb$^{-1}$ at $\sqrt{s} = 546$ GeV [5.12]. The event selection followed very closely the search for ordinary $W \rightarrow e\nu$ and $W \rightarrow \mu\nu$ events, with the additional requirement of observing a photon candidate with $E_T > 10$ GeV. No events that are consistent with the production and decay of a massive $(e^+\nu_e\gamma)$ or $(\mu^+\nu_\mu\gamma)$ state have been found. Using the electron channel, an upper limit on the production of an excited state of charged leptons was derived and expressed in terms of the ratio $R$:

\[
R \equiv \frac{\text{BR}(W \rightarrow e^+\nu_e \rightarrow e\gamma\nu_e) / BR(W \rightarrow e\nu_e)}. \tag{5.8}
\]
Table 5.1
Properties of the $\ell^+ \ell^- \gamma$ events (masses in GeV/c$^2$)

<table>
<thead>
<tr>
<th></th>
<th>e$^+e^-\gamma$</th>
<th>$\mu^+\mu^-\gamma$</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>UA1$^a$</td>
<td>UA2$^a$</td>
</tr>
<tr>
<td>$E_\gamma$ (GeV)</td>
<td>42.9 ± 1.3</td>
<td>24.4 ± 1.4</td>
</tr>
<tr>
<td>$\Delta\alpha(\ell\gamma)$ (°)</td>
<td>11.5</td>
<td>31.4</td>
</tr>
<tr>
<td>$m(\ell\xi)$</td>
<td>42.9 ± 3.4</td>
<td>50.3 ± 1.7</td>
</tr>
<tr>
<td>$m(\ell\delta\gamma)$</td>
<td>104.0 ± 2.2</td>
<td>90.6 ± 2.1</td>
</tr>
<tr>
<td>$m(\ell\gamma)$</td>
<td>3.9 ± 0.3</td>
<td>9.0 ± 0.4</td>
</tr>
<tr>
<td>$m(\ell\gamma)$</td>
<td>94.7 ± 1.9</td>
<td>74.8 ± 2.3</td>
</tr>
</tbody>
</table>

a) 1983 data; b) 1985 data.

Fig. 5.2 Dalitz plots for $\ell^+ \ell^- \gamma$ events [5.10]:

a) for e$^+e^-\gamma$ data (• for UA1 and ■ for UA2);
b) for Standard Model $p\bar{p} \rightarrow \ell^+ \ell^- \gamma + X$;
c) for $Z \rightarrow e^+e^- + e\bar{e}^*$ with $m_{e^*} = 80$ GeV;
d) for $Z \rightarrow Z^*\gamma$ or $\gamma^*\gamma$.  

41
Fig. 5.3 Upper limit (90% CL) for $R$ as a function of $m_{e^*}$ for the process $W \rightarrow e^*\nu_e \rightarrow e\gamma\nu$ obtained by UA1 [5.12].

Figure 5.3 shows the 90% CL upper limit on $R$ as a function of the $e^*$ mass obtained from this analysis.

The UA2 Collaboration has reported results of a search for excited electrons using the full data sample, corresponding to an integrated luminosity of 910 nb$^{-1}$ [3.15]. Requiring an electron candidate with $p_T^e > 11$ GeV/c, a photon candidate with $p_T^\gamma > 10$ GeV/c and $p_T^\gamma > (m_{W} - M^*)/4$ (where $M^*$ is the electron-photon mass), no event was selected. It is expected that more than 95% of the events coming from $W \rightarrow e^*\nu \rightarrow e\gamma\nu$ would survive the above cuts for $M^* \gtrsim 20$ GeV/c$^2$. From a Monte Carlo simulation, the ratio $R$ (Eq. (5.8)) as a function of $M^*$ was extracted. The 95% confidence region in the coupling-strength $\lambda$--$M^*$ plane, for which the excited electron is excluded, is shown in Fig. 5.4. The UA2 result is compared with the most recent limits from $e^+e^-$ data [5.13].

Fig. 5.4 Limits (95% CL) on excited electrons obtained by UA2 [3.15] in the coupling-strength $\lambda$--$M^*$ plane ($M^*$ is the excited electron mass). Also shown are limits from $e^+e^-$ experiments.
6. OTHER NEW PARTICLES

There are many other species of new particles which have been proposed and could be searched for at the CERN p\bar{p} Collider. Here we will limit ourselves to three examples where we are aware of specific experimental searches: leptoquarks, monopoles, and free quarks.

6.1 Leptoquarks

These are strongly interacting colour-triplet particles with couplings to lepton-quark combinations. They appear in several different theoretical frameworks, notably superstring-inspired models [6.1] and some scenarios for compositeness of quarks, leptons [6.2], W, and Z detectable at a low-energy scale. In most such models, any light leptoquark bosons would have spin zero. The superstring-inspired models favour charge |q| = \frac{1}{3} [6.3] whilst the composite models favour charge |q| = \frac{2}{3}. If the leptoquarks have generation-conserving couplings, as is favoured by upper limits on flavour-changing neutral currents, we might expect the following decays in the superstring case:

\[
\begin{align*}
\text{LQ}_1 &\rightarrow d\nu_e, u\tau^-, \\
\text{LQ}_2 &\rightarrow s\nu_e, c\mu^-, \\
\text{LQ}_3 &\rightarrow b\nu_e, t\tau^-,
\end{align*}
\]  

(6.1)

or in the composite model,

\[
\begin{align*}
\text{LQ}_1 &\rightarrow u\bar{\nu}_e, d\mu^+, \\
\text{LQ}_2 &\rightarrow c\bar{\nu}_\mu, s\mu^+, \\
\text{LQ}_3 &\rightarrow t\bar{\nu}_\tau, b\tau^+,
\end{align*}
\]  

(6.2)

In either case, the best signatures are offered by the second-generation leptoquark LQ_2, which can yield $\mu^+\mu^-$ + jet events, $\mu$ + jet ( + $E_T$) events, and jet + $E_T$ events.

In UA1 a search for pair-produced leptoquarks, i.e. $p\bar{p} \rightarrow \text{LQ}_2\overline{\text{LQ}}_2 + X$ was performed using each of these three signatures [6.4]. The ISAJET Monte Carlo program was used for simulating the production and decay of leptoquark events in the UA1 detector. The predicted production cross-section as a function of the leptoquark mass is shown in Fig. 6.1.

To study a possible signal in the $\mu^+\mu^- + \text{jet}$ channel, the UA1 dimuon data sample ($p_T > 3 \text{ GeV/c}$, $m_{\mu\mu} > 6 \text{ GeV/c}^2$) has been used [6.5]. In order to obtain the best possible signal-to-background ratio for $\mu^+\mu^- + \text{jet}$ events coming from LQ production, the selection cuts were modified in the following way: i) isolated muons with $p_T > 5 \text{ GeV/c}$, ii) $\Delta\phi(\mu_1, \mu_2) < 170^\circ$ requiring opposite-sign muons, iii) two jets with $E_T > 8 \text{ GeV}$, iv) $m_{\mu^-\text{jet}} > 15 \text{ GeV/c}^2$ for each combination, and v) $m_{LQ} = m_{\text{LQ}}$ to within 40%. One event survived this selection, to be compared with 1.9 ± 0.8 events expected from the $b\bar{b}$, $c\bar{c}$ production, Drell–Yan, and $\Upsilon$
Fig. 6.1 Predicted leptoquark production cross-section at $\sqrt{s} = 630$ GeV as a function of leptoquark mass obtained with the ISAJET Monte Carlo program [6.4].

decays. Comparison with the rates expected from leptoquarks of different masses (139 events for $m_{LQ} = 20$ GeV/c$^2$, 27 for $m_{LQ} = 30$ GeV/c$^2$, 4.5 for $m_{LQ} = 40$ GeV/c$^2$) gave the 90% CL limits shown in Fig. 6.2, in the plane of $m_{LQ}$ versus the branching ratio for $LQ \rightarrow \mu +$ jet.

The second sample explored was the one containing $\mu +$ jet events. In this case we are looking for events $p\bar{p} \rightarrow LQ\bar{LQ} \rightarrow \mu\bar{\nu}c\bar{s}$, where a high-$p_T$ muon and two jets are expected in the final state. Accordingly the selection was: i) $p_T^\mu > 10$ GeV/c$^2$, and ii) two jets with $E_T^j > 12$ GeV, $E_T^{j2} > 8$ GeV. Additional technical cuts have been applied, tailored to the expected leptoquark signature. To obtain the best possible limit on the leptoquark mass, a likelihood function was constructed to distinguish leptoquark events from background events. The distributions of $p_T^\mu$, of $E_T$ of the second jet ($E_T^{j2}$), of $E_T$, and of $\cos \theta_2^*$ were used for this likelihood function$^*$. Figure 6.3 shows the distribution of the likelihood function for the observed events, compared with that expected for the background (normalized to the data) and for a leptoquark of 30 GeV. The region of L > 0 has been used to obtain the limit from the $\mu +$ jet sample which is also shown in Fig. 6.2.

Finally, UA1 used the $E_T$ event sample discussed in subsection 2.2 to look for double $LQ \rightarrow \nu +$ jet decays. The standard selection, i.e. i) $E_T > 15$ GeV/c$^2$ and $N_\mu > 4$, and ii) $E_T^{j2} > 12$ GeV, was used, together with iii) $E_T^{j1} < 40$ GeV and iv) a tighter cut on $L_T$ to eliminate events with $\tau$ characteristics. This selection left 7 events, to be compared with 10.5 ± 2.9 ± 0.8 expected from standard physics processes and jet fluctuations, and 13 (14) (2.7) events from a leptoquark of mass 20 (30) (40) GeV/c$^2$ decaying entirely into $\nu +$ jet. The deduced limit on the LQ mass is plotted in Fig. 6.2. As can be seen from the excluded regions obtained from the $E_T +$ jet sample, leptoquark masses below ~ 25 GeV/c$^2$ cannot be excluded. At low leptoquark masses, the efficiency for selecting events is very small owing to

*) $\theta_2^*$ is the angle between the antiproton beam direction and the axis of jet 2 in the centre of mass of the system consisting of the muon, jet 1, jet 2, and $E_T$. 

44
the low trigger efficiency. In addition, the 8% systematic error on the jet energy scale prevents the placing of a limit below \( \sim 25 \text{ GeV/c}^2 \). Combining the results of the searches for a second-generation leptoquark in the three different channels discussed, and assuming \( \text{BR}(LQ \to \mu^+ s) + \text{BR}(LQ \to c \bar{p}) = 1 \), led to the bound \( \text{Eqn 6.4} \)

\[
m_{LQ} > 33 \text{ GeV/c}^2, \tag{6.3}
\]

except for a small window where \( 21 \text{ GeV/c}^2 < m_{LQ} < 25 \text{ GeV/c}^2 \) and the branching ratio of \( LQ \to \mu^+ s \) is less than 10%.

### 6.2 Monopoles

It was pointed out many years ago by Dirac [6.6] that magnetic monopoles could be incorporated into QED if their charges \( g \) obeyed a quantization condition:

\[
g \cdot e = 2\pi \cdot n \quad (n \text{ is an integer}), \tag{6.4}
\]

However, the mass of a Dirac monopole was not predictable. More recently, 't Hooft and Polyakov [6.7] have shown that a magnetic monopole is a generic feature of Yang–Mills gauge theory, appearing when a simple non-Abelian gauge group is broken spontaneously to the \( \text{U}(1) \) electromagnetic subgroup. In this case the mass of the monopole is
\[ m_M = O(1) \times m_V/\alpha, \]  
\[ (6.5) \]

where \( m_V \) is the mass acquired by a gauge boson during this spontaneous symmetry breaking. No such monopole appears in the Standard Model, because the gauge group \( SU(2)_L \times U(1)_Y \) is not simple. In many unified theories, the value of \( m_V \) appearing in Eq. (6.5) is much larger than \( m_W \). Even with \( m_W \) in Eq. (6.5), the mass of a ’t Hooft–Polyakov monopole would be beyond the reach of any present accelerator. Therefore one can only search for more general Dirac monopoles at the Collider.

The UA3 experiment has been the only dedicated experiment to search for monopoles. The data-taking period was during the first Collider run operating at \( \sqrt{s} = 546 \text{ GeV} \) [6.8]. The experiment aims at identifying the monopoles by their high ionization rate:

\[ \frac{dE}{dx} \approx \frac{N^2 \beta^2}{(g_D/e)^2}, \]  
\[ (6.6) \]

where \( dE/dx_m \) is the ionization rate for a minimum-ionizing particle, \( \beta \) is the monopole velocity, and \( g_D = \hbar c/2e \).

Plastic detectors, suitable for searching for monopoles, have been placed inside the vacuum pipe of the Collider around the collision point, around the beam pipe, and around the Central Detector of the UA1 experiment, in order to track the monopoles through the magnetic field. No monopole candidates were found. This experiment established an upper limit on the monopole production cross-section of

\[ \sigma \lesssim 2 \times 10^{-31} \text{ cm}^2 \quad (90\% \text{ CL}), \]  
\[ (6.7) \]

for monopoles weighing up to 150 GeV, as seen in Fig. 6.4. Figure 6.4a corresponds to the assumption that fractional charged quarks are elementary \( (g = 3g_D) \); Fig. 6.4b corresponds to the electron being elementary \( (g = g_D) \).

![Fig. 6.4 Cross-section limits (90\% CL) as a function of the monopole mass obtained by UA3 (6.8). Curves are given for two choices of the magnetic charge \( g_D \).](image-url)
6.3 Free Quarks

It is natural to ask whether fractionally charged quarks may become free in nature. According to the standard wisdom in QCD they are confined, but it could be argued that QCD is not quite correct, or that our understanding of QCD is incomplete. If quarks were to become free, some uncertainty would still surround their interaction cross-section with matter, which would affect the strategy for experimental detection. It has often been assumed that quarks would have small interaction cross-sections, but this has been questioned more recently [6.9].

The UA2 Collaboration has searched for free quarks during the first two running periods of the Collider, using a telescope of scintillation counters to detect charged particles with abnormally low ionization densities. Details of the experimental set-up can be found in Ref. [6.10].

Two triggers were used to select events containing low-ionization-density candidates in the quark counters. The first trigger was set up to search for quarks with charge $\pm \frac{1}{3}$, and required a signal in a selected set of counters [imposing an allowed range on the measured amplitude in terms of minimum-ionizing particle (mip) equivalent] in coincidence with a minimum-bias signal. The number of events collected under this trigger condition corresponds to an integrated luminosity of $14.8 \text{ nb}^{-1}$. The second trigger was used to search for quarks with charge $\pm \frac{2}{3}$ requiring a minimum-bias signal together with a signal in all scintillator counters, with an upper limit imposed on their amplitude in terms of mip. This data set corresponds to $3.6 \text{ nb}^{-1}$ of integrated luminosity. The most probable ionization $I_0$ for an event was evaluated from the pulse heights measured in the different counters, using a maximum likelihood method. No candidate events with $I_0 < 0.7 \text{ mip}$ were found for either trigger mode.

The efficiency for detecting quarks with charges equal to $\pm \frac{1}{3}$ and $\pm \frac{2}{3}$ was evaluated by Monte Carlo assuming a momentum distribution of the form $p_T e^{-Bp_T}$, where $m_T = (p_T^2 + m^2)^{1/2}$ and $B = 5 \text{ GeV}^{-1} \text{ c}^2$. The Monte Carlo generated events were subject to the same selection criteria as those used for the experimental data. Figure 6.5 shows the resulting 90% CL upper limits on the ratio $R$.

![Graph showing the 90% CL upper limits on R](image)

Fig. 6.5 The 90% CL upper limit on $R$, the number of quarks per single charged particle as a function of the quark mass obtained by UA2. The labels $\frac{1}{3}$ and $\frac{2}{3}$ refer to the absolute values of the quark charges. The dashed curves represent earlier UA2 results, see Ref. [6.10].
the number of quarks per single charged particle, as a function of the quark mass $m$ for both $\pm \frac{1}{3}$ and $\pm \frac{2}{3}$ electric charges. These limits, of order $2.8 \times 10^{-6}$ and $5.6 \times 10^{-5}$, respectively, for very light quarks, increase rapidly with increasing quark mass. The limits are valid under the assumption that free quarks have the same interaction length as that of ordinary hadrons.

7. SUMMARY AND PREVIEW

The CERN $p\bar{p}$ Collider has not yet discovered any unexpected new particles. The cleanliness of large-$E_T$ events, and the ease with which the $W^+$ and the $Z^0$ were discovered, disarmed many people who were sceptical of high-energy hadron–hadron colliders. Moreover, it is impressive to see how detectors which were designed with other physics objectives in view have been used to make effective searches for new physics signals which had not been foreseen at the time of construction. Here we have in mind, for instance, physics with a missing-energy signature such as $W \rightarrow \tau \nu$ decays, neutrino counting, and searches for heavy leptons and particles. The searches for the latter types of new particles are the most powerful to date, and it is not the fault of the CERN $p\bar{p}$ Collider that these particles do not lie in the mass range accessible to the first generation of Collider experiments.

We extract the following lessons from these searches and from others such as the search for additional heavy quarks. Missing energy will in future be as important a tool for event analysis as are jet energy measurement and lepton identification, and future collider detectors must be able to combine these three features if they seek to explore all the possible new physics sketched in Table 1.1. One corollary is the need for uniform detection over as large a fraction of the solid angle as possible. Missing-energy searches, in particular, are hampered by holes in the azimuthal and/or polar angle coverage. The fine granularity of the detector and the ability to pick out electrons close to jet axes, in addition to muon identification in jets, are also important.

The upgradings of the UA1 and UA2 detectors for running with ACOL reflect the above lessons. The major development of UA1 will be an improved calorimetry with better resolution and finer granularity. The major improvements of UA2 are an improved tracking for better identification of electrons, and an extension in the range of polar angles covered by calorimetry.

The range of new particle masses accessible to these improved detectors depends on the luminosity eventually achieved with ACOL. On the average, one could expect sensitivity to new particle masses up to about 50% larger than those probed so far. The considerably higher centre-of-mass energy of the FNAL Tevatron Collider should eventually pay off with even larger ranges of accessible new particle masses. These should extend up to about twice the present CERN $p\bar{p}$ Collider limit in many cases.
It will be interesting to see whether the e⁺e⁻ colliders which are to operate in the (100 to 200) GeV centre-of-mass energy range (the SLC and LEP) will find any new particles that are inaccessible to the CERN and FNAL pp colliders, such as a Higgs boson, or whether the greater cleanliness of the SLC and LEP finds its major application in precision experiments. We can in any case be sure that the CERN pp Collider, as well as opening up a new range of energy, has also changed the way physicists perceive different types of accelerators. The first explorer in the next energy range up to 1 TeV will presumably be another hadron–hadron collider.

Acknowledgement

It is a pleasure to thank the Scientific Reports Editing and Text Processing Sections at CERN for their competence and patience in preparing this paper.
REFERENCES

    Salam, A., Proc. 8th Nobel Symposium, Aspenäsgården, 1968, ed. N. Svartholm
    (1986).


       See Ref. [2.28].

[2.31] Bloom, E. and Peck, C., in e+e− Annihilation, New Quarks and Leptons, ed. R.N. Cahn,

[2.32] Ellis, J. et al., in Ref. [2.27].


[2.38] Vysotsky, M., in Ref. [2.29].


[3.1] Recent reviews of SUSY include:


[3.2] Fayet, P., Unification of the fundamental particle interactions, eds. S. Ferrara, J. Ellis


[3.5] Ellis, J. et al., in Ref. [3.2].


   Albajar, C. et al. (UA1 Collaboration), paper in preparation.

B259, 519 (1986).
419 (1987).
[5.1] Peskin, M.E., Proc. Int. Symp. on Lepton and Photon Interactions at High Energies,
therein.
Arnison G. et al. (UA1 Collaboration), Phys. Lett. 126B, 398 (1983) and 147B, 241
[6.1] For reviews, see:


   (1985).