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SEARCH FOR SUPERSYMMETRY AT THE $\bar{p}p$ COLLIDER*

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ABSTRACT

Many models of broken supersymmetry predict the existence of supersymmetric fermions $\chi^{\pm,o}$ with masses less than the $W^{\pm}$ and Z^{o}. Often there are two light neutral fermions χ^{o}, even in models with large gaugino masses. The $W^{\pm}$ have large branching ratios for decays into $\chi^{\pm}+\chi^{o}$, with the $\chi^{\pm}$ subsequently decaying into χ^{o} plus hadrons or leptons. We propose looking at the CERN $\bar{p}p$ collider for $W^{\pm}$ production and decay into supersymmetric fermions, a likely signature being "zen" events with one broadened hadronic jet system recoiling against invisible missing transverse energy.

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The CERN $\bar{p}p$ collider has already started making its anticipated discoveries [1,2]. What other adventures may be in store for our experimental colleagues? One possibility is the discovery of supersymmetry [3]. Broken supersymmetric theories are currently the focus of considerable interest, and strategies have been proposed [4] to look in hadron-hadron collisions for strongly interacting supersymmetric particles such as gluinos and squarks. There should in addition be many color singlet supersymmetric fermions $\chi^{\pm,o}$ coupled to the $W^{\pm}$ and the Z^o, and it has been pointed out [5-7] that in many theories at least one particle of each charge should be lighter than the intermediate vector bosons.

In this paper we explore the phenomenology of such fermions in some detail. We discuss the charged and neutral fermion mass matrices, pointing out that there may be light $\chi^{\pm,o}$ even in theories with large gaugino masses. We delineate the areas of parameter space allowed by present experimental searches and by cosmology [8,9]. We point out that in much of the allowed domain one and often two $W^{\pm} \to \chi^{\pm} + \chi^o$ decay channels are open. The branching ratios are expected to be several percent [6,7], while the forward-backward decay asymmetry is model-dependent. Decays involving the lightest neutral supersymmetric fermion, which is probably predominantly a photino $\tilde{\gamma}$, are likely to have a distinctive signature. The charged fermion $\chi^{\pm}$ would recoil against large missing transverse energy, just like the $e^{\pm}$ in the $W^{\pm}$ events already observed [1,2]. The $\chi^{\pm}$ would then decay into χ^o and a pair of charged and neutral conventional leptons or two collimated hadronic jets. Thus a likely signature would be "zen" events* of the type shown in fig. 0: a broadened hadronic jet system in one hemisphere with invisible transverse energy to balance it. We also consider the possibility of similar events from $Z^o \to \chi^o + \chi^{o\prime}$, $\chi^{o\prime} \to \chi^o + X$ decays, but find that

* This terminology is motivated by the zen koan: "You can make the sound of two hands clapping. Now what is the sound of one hand?" See ref. [10] for important background information.

these neutral zen events are suppressed in the cosmologically allowed domains. There could be other sources of zen events, for example, $W^{\pm}$ decays into heavy leptons or Z^o decays into sneutrinos as discussed in ref. [11].

We consider a minimal susy model with two light doublets of Higgs chiral superfields H_1 and H_2 of weak hypercharge ± 1 respectively [3,5,12-15]. The mass matrices for the charged and neutral susy fermions - gauginos and shiggses - are determined by the Lagrangian terms

$$\mathcal{L} \ni + \epsilon \epsilon_{ij} \tilde{H}_1^i \tilde{H}_2^j - M_2 \tilde{W}_a \tilde{W}_a - M_1 \tilde{B} \tilde{B} \tag{1}$$

where W_a and B denote SU(2) and U(1) gauge superfields respectively, the tildes denote fermionic components and i,j (a) are doublet (triplet) SU(2) indices. The quantities ϵ, M_2 and M_1 are mass parameters that are generally expected to be $O(m_W)$. We shall assume

$$M_1 = \frac{5}{3} \frac{\alpha_1}{\alpha_2} M_2 \tag{2}$$

where $\alpha_i \equiv g_i^2/4\pi$, $i = 1,2,3$ are the gauge coupling constants, which holds to leading order in the renormalization group equations [12] if SU(2) $\times$ U(1) is eventually embedded in a unifying non-Abelian group. When combined with the conventional Higgs-gauge field couplings the full mass matrix for the left-handed charged fermion fields becomes

$$\left(\tilde{W}^+, \tilde{H}_1^+\right) \begin{pmatrix} M_2 & g_2 v_2 \\ g_2 v_1 & -\epsilon \end{pmatrix} \begin{pmatrix} \tilde{W}^- \\ \tilde{H}_2^- \end{pmatrix} \tag{3}$$

where $\langle 0|H^o_{1,2}|0\rangle = v_{1,2} : m_W^2 = g_2^2(v_1^2+v_2^2)/2$. This matrix is diagonalized by rotations through angles $\theta_{\pm}$ among the positively and negatively charged fields respectively, where

$$tan\theta_{\pm} = \frac{b_{\pm} + \sqrt{b_{\pm}^2 + 4a_{\pm}^2}}{2a_{\pm}} \tag{4a}$$

with

$$a_{\pm} = \begin{cases} M_2 g_2 v_1 - \epsilon g_2 v_2 \\ M_2 g_2 v_2 - \epsilon g_2 v_1 \end{cases} \quad ; \quad b_{\pm} = M_2^2 - \epsilon^2 \pm g_2^2 (v_2^2 - v_1^2) \quad . \tag{4b}$$

The charged fermion masses are

$$\begin{aligned} m_1 &= M_2\ cos\theta_+ cos\theta_- - g_2 v_2\ cos\theta_+ sin\theta_- - g_2 v_1\ sin\theta_+ cos\theta_- - \epsilon\ sin\theta_+ sin\theta_- \quad , \\ m_2 &= M_2\ sin\theta_+ sin\theta_- + g_2 v_2\ sin\theta_+ cos\theta_- + g_2 v_1\ cos\theta_+ sin\theta_- - \epsilon\ cos\theta_+ cos\theta_- \quad . \end{aligned} \tag{5}$$

Note that in the limit M_2, $\epsilon \to 0$ the charged mass eigenstates become the Dirac fermions $(\tilde{H}_2^-, \bar{\tilde{W}}^+)$ and $(\tilde{W}^-, \bar{\tilde{H}}_1^+)$ (swiggses) with masses

$$g_2\ v_2 \quad , \quad g_2\ v_1 \tag{6a}$$

respectively, while in the limit $M_2 \to \infty$, $\epsilon \to 0$ the eigenmasses become

$$M_2 \quad , \quad \frac{g_2^2\ v_1\ v_2}{M_2} \quad . \tag{6b}$$

Figure 1 displays the mass of the lightest charged fermion $\chi^{\pm}$ for ranges of the unknown parameters (ϵ, M_2, v_1/v_2): since H_1 gives masses to the charge 2/3 quarks and $m_c \gg m_s$, $m_t \gg m_b$ it may well be that $v_1 \geq v_2$. We see that $m_{\chi^{\pm}} \leq m_{W^{\pm}}$ in most of the range of parameters except possibly if both M_2 and ϵ are much greater than M_W.

There are four neutral susy fermions which mix, namely $\tilde{W}^3$, $\tilde{B}^o$, $\tilde{H}_1^o$ and $\tilde{H}_2^o$. Their mixing matrix * is

$$\left(\tilde{W}^3, \tilde{B}^o, \tilde{H}_1^o, \tilde{H}_2^o\right) \begin{pmatrix} M_2 & 0 & -\frac{g_2 v_1}{\sqrt{2}} & \frac{g_2 v_2}{\sqrt{2}} \\ 0 & \frac{5}{3}\frac{\alpha_1}{\alpha_2} M_2 & \frac{g_1 v_1}{\sqrt{2}} & -\frac{g_1 v_2}{\sqrt{2}} \\ -\frac{g_2 v_1}{\sqrt{2}} & \frac{g_1 v_1}{\sqrt{2}} & 0 & \epsilon \\ \frac{g_2 v_2}{\sqrt{2}} & -\frac{g_1 v_1}{\sqrt{2}} & \epsilon & 0 \end{pmatrix} \begin{pmatrix} \tilde{W}^3 \\ \tilde{B}^o \\ \tilde{H}_1^o \\ \tilde{H}_2^o \end{pmatrix} \tag{7}$$

* Here we correct errors in ref. [5].

which in terms of the convenient combinations

$$\tilde{A}^o \equiv \frac{v_1 \tilde{H}_1^o - v_2 \tilde{H}_2^o}{v} \quad , \quad \tilde{S}^o \equiv \frac{v_2 \tilde{H}_1^o + v_1 \tilde{H}_2^o}{v} \tag{8}$$

(where we have introduced $v \equiv \sqrt{v_1^2 + v_2^2}$) becomes

$$\left(\tilde{W}^3, \tilde{B}^o, \tilde{A}^o, \tilde{S}^o\right) \begin{pmatrix} M_2 & 0 & -\frac{g_2 v}{v_2} & 0 \\ 0 & \frac{5}{3}\frac{\alpha_1}{\alpha_2} M_2 & \frac{g_1 v}{\sqrt{2}} & 0 \\ -\frac{g_2 v}{\sqrt{2}} & \frac{g_1 v}{\sqrt{2}} & -\frac{2 v_1 v_2}{v}\epsilon & \frac{v_1^2 - v_2^2}{v^2}\epsilon \\ 0 & 0 & \frac{v_1^2 - v_2^2}{v^2}\epsilon & \frac{2 v_1 v_2}{v^2}\epsilon \end{pmatrix} \begin{pmatrix} \tilde{W}^3 \\ \tilde{B}^o \\ \tilde{A}^o \\ \tilde{S}^o \end{pmatrix} . \tag{9}$$

The matrix (9) can be diagonalized by an orthogonal rotation whose general form is complicated. In the limit M_2, $\epsilon \to 0$ the mass eigenstates become

$$\tilde{\gamma} \equiv \frac{g_1 \tilde{W}^3 + g_2 \tilde{B}^o}{\sqrt{g_1^2 + g_2^2}} \qquad : m_{\tilde{\gamma}} \approx \frac{8}{3} \frac{g_1^2}{g_1^2 + g_2^2} M_2 \tag{10a}$$

$$\tilde{Z}_\pm^o \equiv \frac{g_1 \tilde{B}^o - g_2 \tilde{W}^3 \pm \sqrt{g_1^2 + g_2^2}\ \tilde{A}^o}{\sqrt{2(g_1^2 + g_2^2)}} : m_{\tilde{Z}_\pm^o} \approx m_{Z^o} = \sqrt{\frac{g_1^2 + g_2^2}{2}}\ v \tag{10b}$$

$$\tilde{S}^o \qquad : m_{\tilde{S}^o} = \frac{2 v_1 v_2}{v^2}\ \epsilon \tag{10c}$$

The Higgs combinations $\tilde{A}^o$ and $\tilde{S}^o$ are approximate mass eigenstates in the limit $M_2 \to \infty$ but ϵ small. Figure 1 displays the two lightest neutral fermion masses. We see that in general the lightest neutral fermion is lighter than the lightest charged fermion. This is favored by cosmology, since the lightest supersymmetric particle is essentially stable, and if it were charged it would dissipate and condense in conventional matter with a density far above experimental upper limits on stable exotic relics from the big bang [16]. We also see from fig. 1 that in general there are two light neutral fermions, one of which is predominantly the photino when M_2, $\epsilon \ll m_W$, while the other is mainly a shiggs as seen in eq. (10c).

There are significant constraints on the parameters M_2 and ϵ which come from particle physics [17] and cosmology [8,9]. No new charged fermion has been seen [17] with a mass less than about 20 GeV and we shall take this as a lower limit on $m_{\chi^\pm}$, though it has been argued [15] that a light charged supersymmetric fermion might have escaped detection. Cosmology imposes an upper limit of at most 2×10^{-29} gm/cc on the possible density * of stable heavy neutral fermions [16]. Their density is reduced to acceptable levels only if they annihilate sufficiently efficiently, and it has recently been pointed out [8] that the annihilation of Majorana fermions like ours is strongly suppressed by P-wave phase space at low temperatures. The effective interaction for annihilation into Dirac fermions $\bar{f}\,f$ has the general form

$$\mathcal{L} = \bar{\psi}\ \gamma^\mu\ \gamma_5\ \psi\ \bar{f}\ \gamma_\mu (AP_L + BP_R) f \tag{11}$$

where A, B receive contributions from intermediate Z^o and sfermion $\tilde{f}$ exchange in the s and t channels respectively. If we define an arbitrary neutral Majorana fermion

$$\psi \equiv \begin{pmatrix} \alpha\,\tilde{W}^3 + \beta\,\tilde{B} + \gamma\,\tilde{H}_1^o + \delta\,\tilde{H}_2^o \\ \alpha\,\bar{\tilde{W}}^3 + \beta\,\bar{\tilde{B}} + \gamma\,\bar{\tilde{H}}_1^o + \delta\,\bar{\tilde{H}}_2^o \end{pmatrix} \tag{12}$$

and use a convention where $Q_f = T_f^3 + (1/2)Y_f$ then the Z^o and sfermion exchange contributions to A and B are

*This bound comes from the overall density of the Universe and is very conservative. One can argue that massive neutral fermions probably condense into galaxies in which case a more stringent limit coming from missing galactic matter could be applied. See ref. [9] for more details.

$$A_f^Z = (\gamma^2 - \delta^2)\ \frac{g_1\ sin\theta_W + g_2\ cos\theta_W}{4M_Z^2}\Big[\frac{1}{2}\ Y_{f_L} g_1\ sin\theta_W - T^3_{f_L} g_2\ cos\theta_W\Big] \tag{13a}$$

$$B_f^Z = (\gamma^2 - \delta^2)\ \frac{g_1\ sin\theta_W + g_2\ cos\theta_W}{8M_Z^2}\ Y_{f_R} g_1\ sin\theta_W \tag{13b}$$

$$A_f^{\tilde{f}} = \frac{\left(T^3_{f_L}\alpha g_2 + \frac{1}{2}\ Y_{f_L}\beta g_1\right)^2}{2m^2_{\tilde{f}_L}} \tag{13c}$$

$$B_f^{\tilde{f}} = -\frac{\left(\frac{1}{2}\ Y_{f_R}\beta g_1\right)^2}{2m^2_{\tilde{f}_R}} \tag{13d}$$

where additional small contributions to $A_f^{\tilde{f}}$, $B_f^{\tilde{f}}$ involving shiggs couplings proportional to fermion masses are omitted for simplicity of presentation.* The gaugino couplings (13c,d) can be larger than the shiggs couplings (13a,b) if $m_{\tilde{f}} < m_Z$ (as allowed by existing constraints [17]), and thus the gauginos potentially annihilate more efficiently. We find that to be consistent with cosmology, the lightest susy particle should not be predominantly a shiggs, but should be a photino or contain gaugino components. Figure 2 shows the ranges of M_2 and ϵ which are consistent with cosmology if $m_{\tilde{q},\tilde{\ell}} \approx$ 20 GeV which is the most favorable case for annihilation and hence yields the most conservative bounds. We refer the interested reader elsewhere [9] for a more complete study of the cosmological constraints on susy particles.

There are substantial regions of the experimentally and cosmologically allowed domains in fig. 2 for which the decays $W^{\pm} \to \chi^{\pm} + \chi^{o}$ are kinematically accessible [6].

* These additional shiggs contributions to the sfermion exchange diagram cannot in general be cast into an effective interaction of the form (11): while the contribution quadratic in the shiggs couplings can, the shiggs-gaugino interference term cannot. We have included the shiggs-shiggs contribution in computing the bounds in fig. 2, and found that it becomes important for the annihilation of neutral shiggses if $v_1 = v_2$, since the Z^o contribution to the annihilation cross section vanishes in this limit ($\gamma^2 = \delta^2$).

Indeed, there are sizeable regions with two available decay modes involving predominantly photino and shiggs states respectively, which are also shown in fig. 2. We have studied the branching ratios and forward-backward asymmetries for $W^{\pm}$ decay into these modes for general values of the parameters ϵ, M_2 and v_1/v_2. If we define a neutral mass eigenstate as in (11) and take the charged Dirac eigenstates from eq. (4):

$$\chi_1^- = \begin{pmatrix} \cos\theta_- \tilde{W}^- - \sin\theta_- \tilde{H}_2^- \\ \cos\theta_+ \bar{\tilde{W}}^+ - \sin\theta_+ \bar{\tilde{H}}_1^+ \end{pmatrix} \quad , \quad \chi_2^- = \begin{pmatrix} \sin\theta_- \tilde{W}^- + \cos\theta_- \tilde{H}_2^- \\ \sin\theta_+ \bar{\tilde{W}}^+ + \cos\theta_+ \bar{\tilde{H}}_1^+ \end{pmatrix} \tag{14}$$

then, for example, the left- and right-handed couplings of the W^- to the $\chi_1^- \chi^o$ combination are (in units of $g_2/\sqrt{2}$)*

$$\begin{aligned} g_L &= \sqrt{2}\ \cos\theta_- \alpha - \sin\theta_- \delta \\ g_R &= -\sqrt{2}\ \cos\theta_+ \alpha + \sin\theta_+ \gamma \quad . \end{aligned} \tag{15}$$

The branching ratio for $W^{\pm} \to \chi^{\pm} + \chi^o$ is simply

$$R \equiv \frac{B(\chi^{\pm}\chi^o)}{B(e^{\pm}\nu)} = \frac{6p}{M_W^3}\left[(g_L^2 + g_R^2)\left(E_{\pm}E_0 + \frac{p^2}{3}\right) + 2g_L g_R\ m_{\pm} m_0\right] \tag{16a}$$

where

$$p = \frac{1}{2}\sqrt{M_W^2 - 2m_{\pm}^2 - 2m_0^2 + (m_{\pm}^2 - m_0^2)^2 / M_W^2} \tag{16b}$$

is a final state three-momentum.

Defining "forward" to be when a negative particle emerges in the direction of the proton beam, the forward-backward asymmetry is

$$A \equiv \frac{\int_0^1 d\sigma - \int_{-1}^0 d\sigma}{\int_0^1 d\sigma + \int_{-1}^0 d\sigma} = \frac{p m_W (g_L^2 - g_R^2)}{2\left[(g_L^2 + g_R^2)(E_{\pm}E_0 + p^2/3) + 2g_L g_R\ m_{\pm} m_0\right]} \tag{17}$$

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*If a charged or neutral eigenstate corresponds to a negative eigenvalue of the mass matrices (3) or (9), it is necessary to introduce a relative minus sign between the top (left) and bottom (right) Weyl components of the Dirac spinor to obtain the physical mass eigenstate. This would flip the relative sign of g_L and g_R in (15).

so that the electron asymmetry defined by (17) is expected to be +3/4. Figure 3 shows how R and A vary for $W^{\pm}$ decays into $\chi^{\pm}$ + the lightest χ^{o}, generally predominantly a $\tilde{\gamma}$. The rates for $W^{\pm} \rightarrow \chi^{\pm} + \chi^{o\prime}$ are not shown: they are generally larger because the SU(2) gauge coupling to the shiggs is larger than the electromagnetic coupling of the photino. We have chosen to emphasize decays into the lightest neutral fermion because they are more experimentally accessible, as we will see in a moment. We see that the $W^{\pm} \rightarrow \chi^{\pm} + \chi^{o}$ rates are a substantial fraction R of the $W^{\pm} \rightarrow e^{\pm} + \nu$ branching ratio which is expected to be about 8% in the standard model. The susy fermions have an essentially null forward-backward asymmetry if $v_1 \approx v_2$, but may have a substantial negative asymmetry if $v_1 \gg v_2$. These results can be related intuitively to the $\tilde{W}$ and $\tilde{H}$ contents of the charged mass eigenstates. They contrast with the asymmetry of +3/4 expected for the electron or for a further sequential heavy lepton, and can in principle be used to distinguish between them and susy particles if one can construct a measure of the $\chi^{\pm}$ charge, for example by weighting suitably the charges of its decay products as functions of their momenta [18].

The charged susy fermions $\chi^{\pm}$ are expected to decay via $W^{\pm}$ or sfermion exchange into some $\chi^{o} + (e\nu, \mu\nu, \tau\nu$, or $q\,\bar{q})$. The lifetime is not expected to be long enough for the $\chi^{\pm}$ decay path to be observable. The final state particle distributions should resemble those in conventional heavy lepton decay, though the Michel parameter will not in general correspond to pure $(V - A)$ or $(V + A)$ interactions. A heavier $\chi^{o\prime}$, if produced, would decay into the lighter $\chi^{o} + (\nu\,\bar{\nu}, \ell^{+}\ell^{-}$, or $q\,\bar{q})$. This can occur via sfermion exchange, or via Z^{o} exchange since the Z^{o} couplings to χ^{o} and $\chi^{o\prime}$ are not diagonal (and in fact become purely off-diagonal if $v_1 = v_2$). The Z^{o} contribution requires a shiggs component in the predominantly photino χ^{o} eigenstate, but otherwise gives a contribution to the $\chi^{o\prime}$ rate which is typical of a heavy lepton. A potentially larger contribution to the $\chi^{o\prime}$ decay rate comes from sfermion exchange, which occurs either through the gaugino component of $\chi^{o\prime}$ or directly through the shiggs-fermion coupling $\propto m_f/M_W$. Since there is typically ample phase space for $\chi^{o\prime} \to \chi^o + X$ decay, it is unlikely that the $\chi^{o\prime}$ would live long enough for its decay to provide a separated vertex, so we shall not examine this possibility further.

The promising signature to search for experimentally at the $\bar{p}p$ collider is likely to be $W^\pm \to \chi^\pm$ + lightest χ^o, $\chi^\pm \to$ lightest $\chi^o + (\bar{q}\,q)$. Figure 3 shows that these events may occur with a rate close to that for $W^\pm \to e^\pm \nu$ decay. Comparison of figs. 1 and 3 shows that in the interesting region the $\chi^\pm$ typically has a mass O(30) GeV, while the lightest χ^o has a mass O(10) GeV. The $\chi^\pm$ is therefore usually produced relativistically with a transverse energy O(40 to 50) GeV, and with a recoil transverse energy O(30 to 40) GeV. The $\chi^\pm$ decays into a missing neutral with transverse energy O(15) GeV and two hadronic jets with invariant mass O(20) GeV and total transverse energy O(25) GeV. The resulting event signature is shown in fig. 0: two collimated jets coalescing into a broadened hadronic jet on one side of the beam axis with a net missing recoil energy of O(25) GeV on the opposite side. These are what we call "zen events." If the $\chi^\pm$ mass is increased, the two hadronic jets become more splayed out, until in an idealized case the $\chi^\pm$ becomes nonrelativistic in the $W^\pm$ rest frame, and the typical azimuthal angle between the two hadronic jets may approach 120^o. In this case the event structure resembles more closely the form already discussed by other authors [4] in connection with gluino or squark pair-production. It is a general feature of supersymmetric theories that one expects events with large amounts of missing transverse energy-momentum which do not contain leptons. In this respect they resemble heavy lepton production events, and we note parenthetically that a search for zen events can also be interpreted as a search for heavy leptons with masses between the present experimental limit [17] of about 18 GeV and m_W. Ways of distinguishing between heavy leptons and susy fermions include the forward-backward production asymmetry, the Michel decay parameter ρ which must be 3/4 for a conventional sequential heavy lepton but could be 0 for $\chi^{\pm}$ decay thus modifying the final state jet energy distributions, and the possibility of unconventional decay branching ratios due to sfermion exchange diagrams.

We have also studied the decay of the Z^o into light neutral susy fermions. In general, there are three kinematically allowed modes: $Z^o \to \chi^o\chi^o$, $Z^o \to \chi^o\chi^{o\prime}$ and $Z^o \to \chi^{o\prime}\chi^{o\prime}$ where χ^o and $\chi^{o\prime}$ are respectively the lightest and next-to-lightest neutral susy fermion. We have previously observed that cosmological arguments prefer a light photino to a light shiggs. Furthermore, since $Z^o \to \tilde{\gamma}\tilde{\gamma}$ is strongly suppressed relative to $Z^o \to \tilde{H}\tilde{H}$, this implies $\Gamma(Z^o \to \chi^o\chi^o) \ll \Gamma(Z^o \to \chi^o\chi^{o\prime}) \ll \Gamma(Z^o \to \chi^{o\prime}\chi^{o\prime})$. The most likely mode for susy Z^o decay is therefore $Z^o \to \chi^{o\prime}\chi^{o\prime}$, where both of the $\chi^{o\prime}$ decay into χ^o plus quarks or leptons. The experimental signature for such a process is unfortunately less distinctive in $\bar{p}p$ collisions than a bona fide zen event. However $e^+e^- \to \chi^{o\prime}\chi^{o\prime}$ through a virtual Z^o could be looked for at present day as well as forthcoming e^+e^- machines. The cross section for this process is

$$\sigma(e^+e^- \to \chi^{o\prime}\chi^{o\prime}) = (\gamma^2-\delta^2)^2\,\frac{G_F^2 s}{12\pi}\,\frac{\left(g_V^{e2}+g_A^{e2}\right)}{\left(1-s/M_Z^2\right)^2+\Gamma_Z^2/M_Z^2}\left(1-\frac{4m_{\chi^{o\prime}}^2}{s}\right)^{3/2} \tag{18}$$

where $g_V^e = -\frac{1}{2}+2sin^2\theta_W$, $g_A^e = -1/2$, or simply $(\gamma^2-\delta^2)^2$ in units of a conventional neutrino cross-section if $m_{\chi^{o\prime}}^2 \ll s$. Typically we find that $(\gamma^2-\delta^2)^2$ is $O(1)$ in the cosmologically allowed domain, except when $v_1 = v_2$ where $(\gamma^2-\delta^2)$ vanishes identically.

We note that for total center-of-mass energies far below the Z^o resonance, a second production mechanism for neutral susy fermions due to selectron exchange can be important if the selectron is light enough [15,19]. Since the electron-shiggs coupling is negligible, this mechanism requires a significant gaugino component within the final state fermions, so that in general one expects $\sigma_{\tilde{e}}(e^+e^- \to \chi^o\chi^o) \gg \sigma_{\tilde{e}}(e^+e^- \to$

$\chi^o\chi^{o\prime}) \gg \sigma_{\tilde{e}}(e^+e^- \rightarrow \chi^{o\prime}\chi^{o\prime})$. It is quite possible that the $\chi^{o\prime}$ can contain sufficient gaugino components for this mechanism to produce observable $e^+e^- \rightarrow \chi^o\chi^{o\prime}$ zen events at present energies, as discussed in ref. [15].

We conclude that the rate of zen events from $W^\pm$ decay and the magnitude of their missing energy-momentum seems to place them well within the reach of experiments [1,2] with the CERN $\bar{p}p$ collider in the near future. Let us hope our experimental colleagues are lucky enough to make another exciting discovery.

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FIGURE CAPTIONS

0. Zen event signature, with a charged susy fermion $\chi^{\pm}$ decaying into a spray of hadrons on one side of the beam axis denoted by $\otimes$, and transverse energy-momentum balanced by two light neutral supersymmetric fermions χ^{o}.

1. Charged and neutral mass eigenstates for plausible ranges of $|\epsilon|/m_W$ and M_2/m_W, assuming (a) $v_1 = v_2$, $\epsilon > 0$; (b) $v_1 = v_2$, $\epsilon < 0$; (c) $v_1 = 4v_2$, $\epsilon > 0$; (d) $v_1 = 4v_2$, $\epsilon < 0$. Solid lines correspond to the lightest neutral eigenstate; dashed lines to the next-to-lightest neutral eigenstate. Dotted lines denote the lightest charged eigenstate.

2. Domains of parameter space consistent with cosmology (solid) and PEP/PETRA limits (dashed). Also shown are the domains in which one (hatched) and two (cross-hatched) $W^{\pm} \to \chi^{\pm} + \chi^{o}$ decay modes are kinematically allowed. The labels (a) to (d) correspond to those in fig. 1.

3. Rates and forward-backward asymmetries for zen events in the allowed regions of fig. 2. Dotted lines are rates normalized to the $e\nu$ rate. Dashed lines represent forward-backward asymmetries. The labels (a) to (d) correspond to those in figs. 1 and 2.

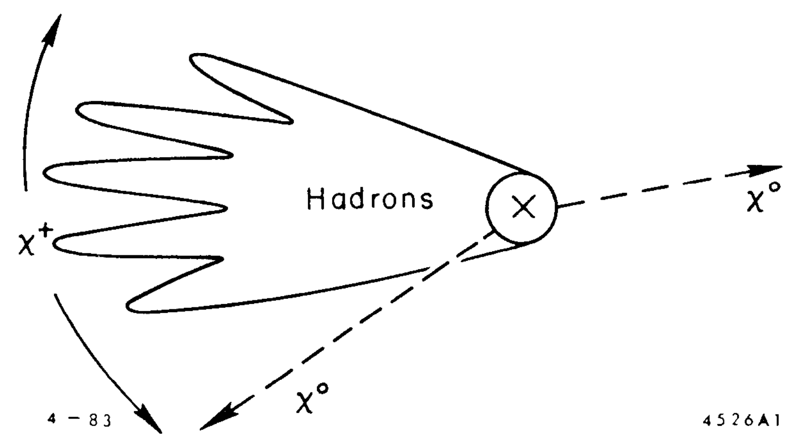


Fig. 0

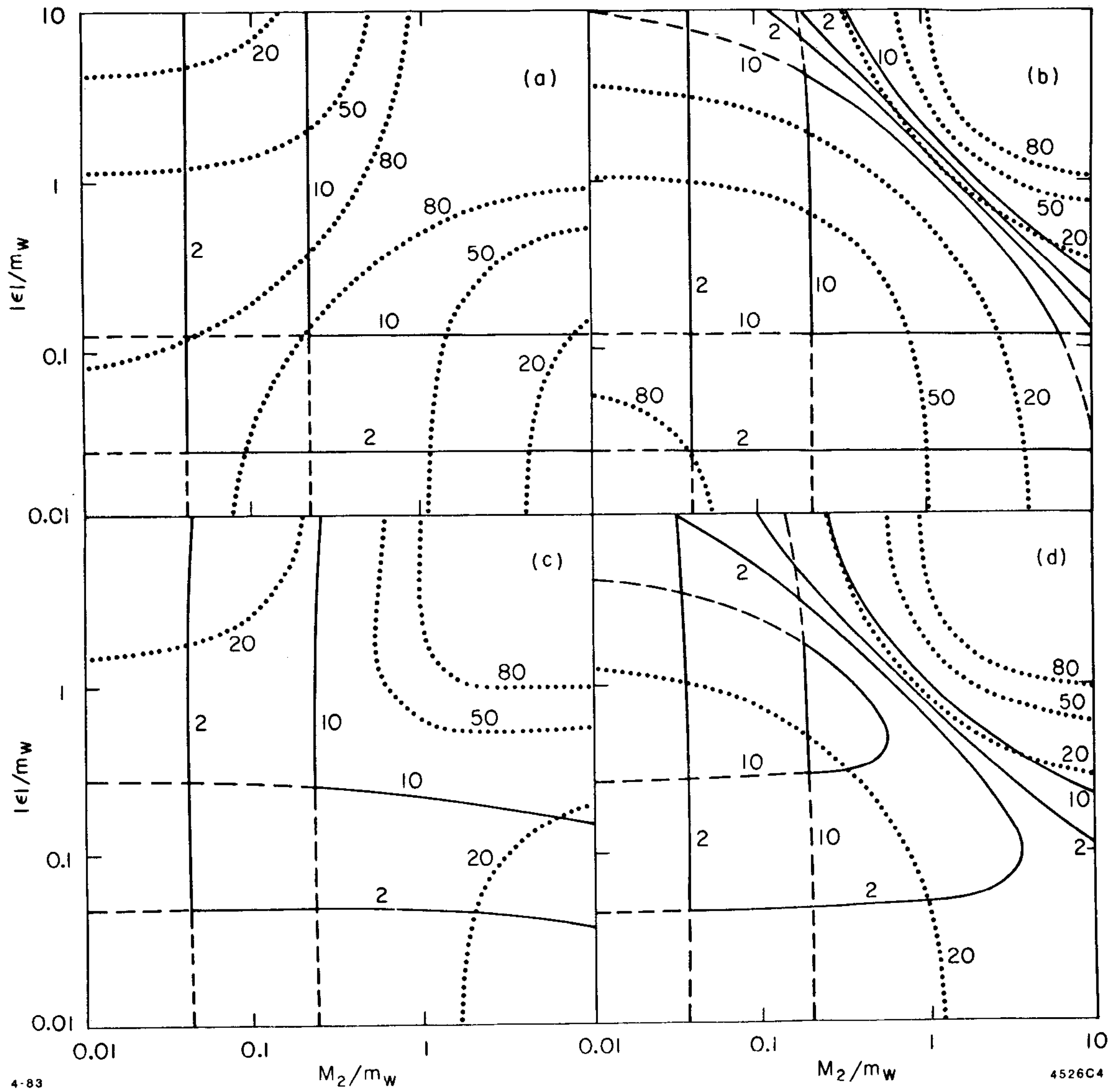
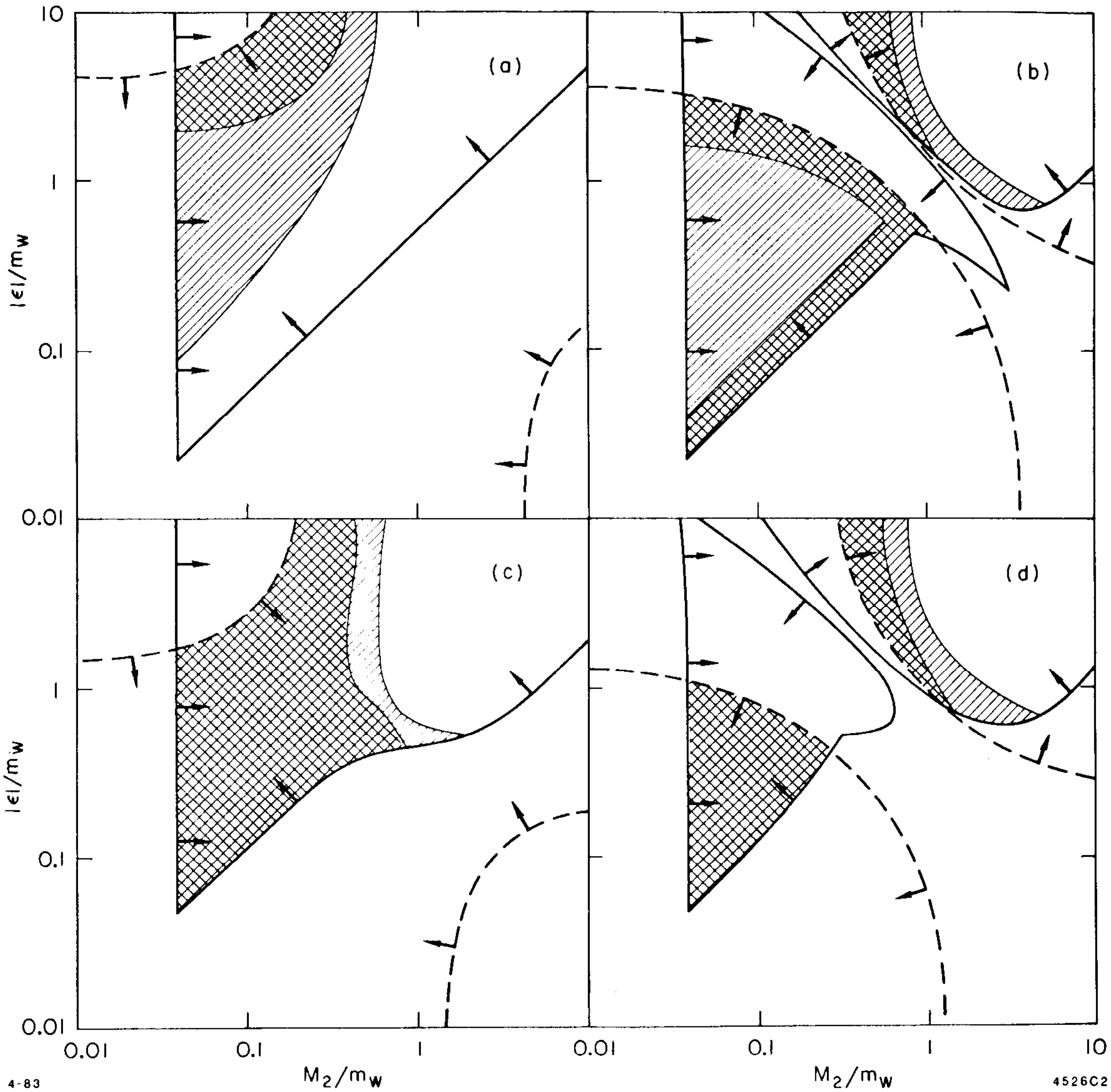


Fig. 1



(a)
(b)
(c)
(d)
10
1
0.1
0.01
$|\epsilon|/m_W$
$|\epsilon|/m_W$
0.01
0.1
1
10
M_2/m_W
M_2/m_W
4-83
4526C2

Fig. 2

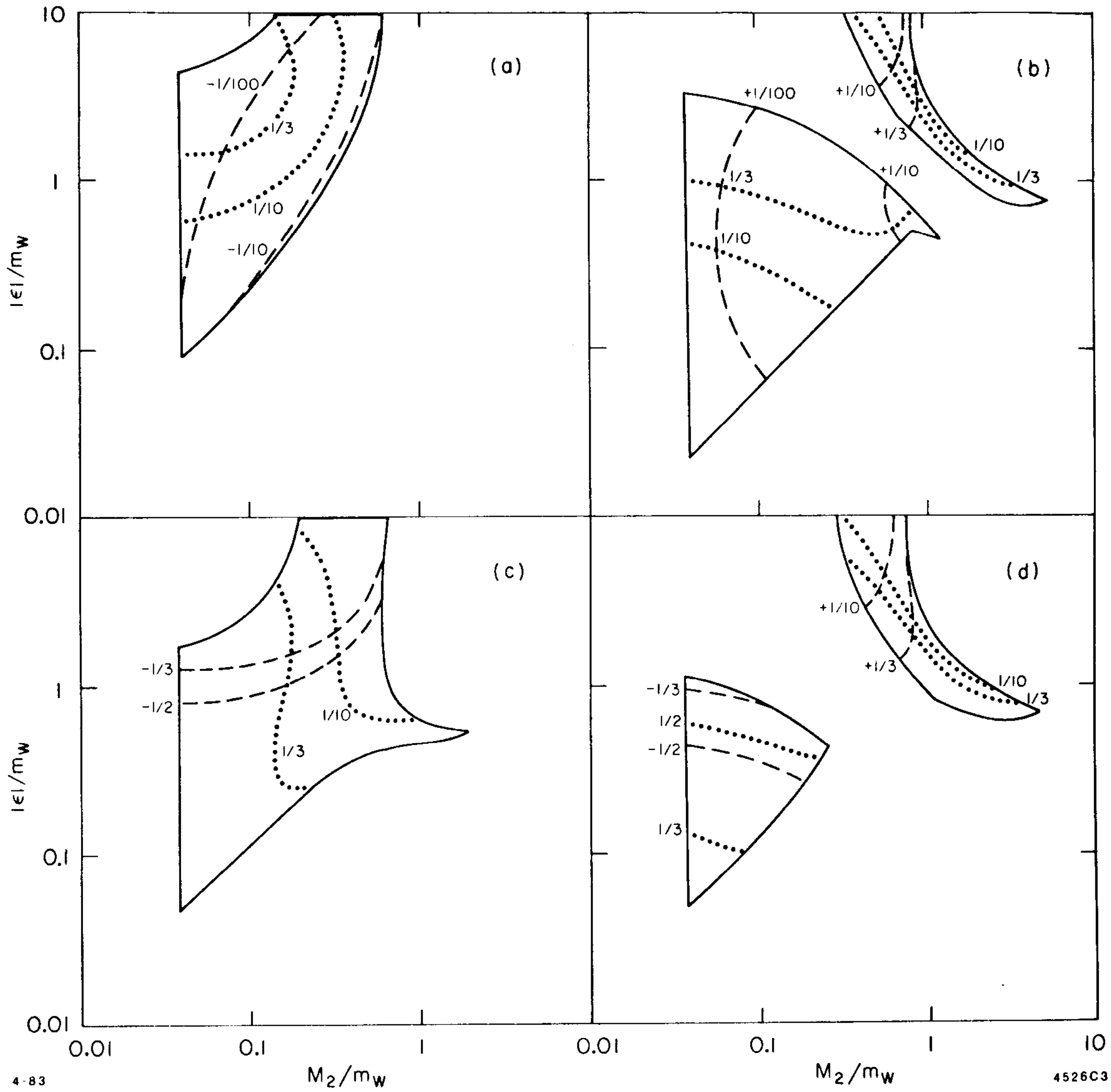

(a)
(b)
(c)
(d)
-1/100
1/3
1/10
-1/10
+1/100
+1/10
+1/3
-1/3
-1/2
1/2
$|\epsilon|/m_W$
M_2/m_W
10
1
0.1
0.01
4-83
4526C3

Fig. 3