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WHAT IF THE HIGGSINO IS THE LIGHTEST SUPERSYMMETRIC PARTICLE?*

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ABSTRACT

A pedagogical introduction to the mixing of neutral gauginos and Higgsinos in supersymmetric models is given. The possibility that the Higgsino (rather than the photino) is the lightest supersymmetric particle is considered and implications for phenomenology are discussed with some emphasis on signatures of supersymmetry in Z^0 decays. Some related aspects of Higgs boson detection in Z^0 decays are mentioned.

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1. Introduction

In most papers which discuss the phenomenological implications of "low energy" supersymmetry,[1–3] it is usually assumed that the photino ($\tilde{\gamma}$) is the lightest supersymmetric particle (LSP). In nearly all supersymmetric models, all particles possess a conserved multiplicative quantum number called R-parity given by $R = (-1)^{3B+L+2J}$ where B is the baryon number, L is the lepton number and J is the spin of the particle.[4] (The possibility that R-parity is broken has been discussed in the literature.[5,6] We shall assume here that R-parity is an exact discrete symmetry.) It is easy to show that all presently observed particles have $R = 1$, whereas their supersymmetric partners would have $R = -1$. Cosmological arguments suggest that the LSP must be both color and electrically neutral.[7] Then, a conserved R-parity implies that the LSP is exactly stable, weakly interacting and will be produced in the decay products of any heavier supersymmetric particle. That is, the LSP behaves like a neutrino and will typically escape collider detectors. Thus, the best experimental signature of supersymmetry is to find evidence for a new neutral weakly interacting particle which is *not* the neutrino.

The assumption that the photino is the LSP implies a very definitive pattern to supersymmetric phenomenology. For completeness, it is important to consider other possibilities. One possible alternative candidate is the scalar-neutrino[8] ($\tilde{\nu}$). If the photino were the second lightest supersymmetric particle, then the photino would be unstable, decaying via $\tilde{\gamma} \to \nu\overline{\tilde{\nu}}$. (This occurs by a one-loop Feynman diagram similar to the decay $\tilde{\nu} \to \nu\tilde{\gamma}$ in the usual scenario.[9]) Because the ν and $\tilde{\nu}$ will escape collider detectors, the $\tilde{\gamma}$ will also be unobserved and the phenomenology will be unchanged as compared with the usual case where the $\tilde{\gamma}$ is the LSP. The other possibility for the LSP is a neutralino (i.e. some mixture of neutral gaugino and Higgsino) which is not a pure photino. In Section 2, we will discuss some general features of neutralino mixing and investigate a particular special case where a pure Higgsino ($\widetilde{H}^0$) is the LSP. In Section 3, we discuss the phenomenology of a supersymmetric model with the $\widetilde{H}^0$ as the LSP. Our conclusions are summarized in Section 4.

2. Neutralino Mixing

Let us begin by making a list of all neutral color-singlet states with zero lepton number in the minimal supersymmetric extension of the Standard Model. (The last restriction has been imposed to eliminate neutrinos and their superpartners from the discussion.) One particular noteworthy feature of this model is the existence of two complex Higgs doublets which are necessary (in a supersymmetric model) to give mass to both up and down type quarks.[10,11] The complete list of particles with the above properties in this model is:

$$B, W_3 \text{ (neutral gauge bosons)} \tag{1a}$$

$$H_1, H_2 \text{ (neutral complex Higgs bosons)} \tag{1b}$$

$$\widetilde{B}, \widetilde{W}_3 \text{ (neutral gauginos)} \tag{2a}$$

$$\widetilde{H}_1, \widetilde{H}_2 \text{ (neutral Higgsinos)} \tag{2b}$$

We have listed the "interaction" eigenstates, i.e. the fields that have definite $SU(2) \times U(1)$ quantum numbers. The physically observable particles are the "mass" eigenstates which are obtained by diagonalizing the appropriate mass matrices. For example, as every student knows, the physical neutral gauge bosons are obtained by diagonalizing the 2×2 matrix:

$$\begin{pmatrix} g'^2 & -gg' \\ -gg' & g^2 \end{pmatrix} \tag{3}$$

The resulting states are the Z^0 and γ:

$$Z^0 = W_3 \cos\theta_W - B \sin\theta_W \tag{4a}$$

$$\gamma = W_3 \sin\theta_W + B \cos\theta_W \tag{4b}$$

where $\tan\theta_W = g'/g$. In the Higgs boson sector, the diagonalization procedure is more complicated but very straightforward. The mass matrices can be found in Ref. 12; after diagonalization, the resulting states are:

$$H_1^0 = \sqrt{2}\,[(Re\ H_1 - v_1)\cos\alpha + (Re\,H_2 - v_2)\sin\alpha] \tag{5a}$$

$$H_2^0 = \sqrt{2}\,[-(Re\,H_1 - v_1)\sin\alpha + (Re\,H_2 - v_2)\cos\alpha] \tag{5b}$$

$$H_3^0 = \sqrt{2}\,[Im\,H_1 \sin\beta + Im\,H_2 \cos\beta] \tag{5c}$$

and

$$G^0 = \sqrt{2}\,[Im\,H_1 \cos\beta - Im\,H_2 \sin\beta] \tag{6}$$

where

$$\tan\beta = \frac{v_2}{v_1} \tag{7}$$

and the angle α depends on coefficients which appear in the Higgs potential. The vacuum expectation values of the neutral Higgs bosons are denoted by $\langle H_i \rangle \equiv v_i$; in this model, we may choose a phase convention such that v_1 and v_2 are both real and positive. This implies that $0 \leq \beta \leq \frac{\pi}{2}$. Note the counting of degrees of freedom here: we started off with two *complex* neutral fields (Eq. (1b)) and ended up with four *real* neutral fields (Eqs. (5), (6)). If one computes the masses of the scalar fields given by Eqs. (5) and (6), one finds that the mass of G^0 is zero. In fact, one can show that this field is the Goldstone boson which is absorbed ("eaten") by the Z^0 when the Z^0 gains mass via the Higgs mechanism. Thus, at the end, three physical neutral Higgs bosons (given by Eq. (5)) remain. By examining the couplings of these physical Higgs bosons to fermions, one finds that H_3^0 is a pseudoscalar and H_1^0 and H_2^0 are scalars. The masses of these neutral Higgs bosons are not independent due to the relations of various parameters of the Higgs potential which are imposed by supersymmetry. One finds after diagonalizing the Higgs boson mass matrix that:[11-13]

$$m^2_{H_1^0,H_2^0} = \frac{1}{2}\left[m^2_{H_3^0} + m_Z^2 \pm \sqrt{\left(m^2_{H_3^0} + m_Z^2\right)^2 - 4m_Z^2 m^2_{H_3^0}\cos^2 2\beta}\right] . \tag{8}$$

By convention we take H_2^0 to be the lighter of the two boson masses in Eq. (8). Note that Eq. (8) implies that $m_{H_1^0} \geq m_Z$ and $m_{H_2^0} \leq m_Z$. The mixing angle α (defined in Eq. (5)) is given by:

$$\sin 2\alpha = -\sin 2\beta \left(\frac{m^2_{H_1^0} + m^2_{H_2^0}}{m^2_{H_1^0} - m^2_{H_2^0}}\right) . \tag{9}$$

Using the conventions stated above (where $0 \leq \beta \leq \frac{\pi}{2}$), it follows that we may choose $-\frac{\pi}{2} \leq \alpha \leq 0$ (since by Eq. (9), $\sin 2\alpha$ is negative). This completes the analysis of the particles listed in Eq. (1).

We now repeat the analysis with the supersymmetric partners of the gauge and Higgs bosons (see Eq. (2)). These particles are all fermions with identical conserved quantum numbers (e.g. they are color and electrically neutral); hence they can all mix.[14] We employ the term "neutralinos" to describe the mass-eigenstates one obtains by diagonalizing the mass matrix, as these states can be complicated linear combinations of gauginos and Higgsinos. In the minimal supersymmetric model the neutralino mass matrix is given by:[15]

$$\begin{pmatrix} M_1 & 0 & -m_Z \sin\theta_W \cos\beta & m_Z \sin\theta_W \sin\beta \\ 0 & M_2 & m_Z \cos\theta_W \cos\beta & -m_Z \cos\theta_W \sin\beta \\ -m_Z \sin\theta_W \cos\beta & m_Z \cos\theta_W \cos\beta & 0 & -\mu \\ m_Z \sin\theta_W \sin\beta & -m_Z \cos\theta_W \sin\beta & -\mu & 0 \end{pmatrix} \tag{10}$$

where the matrix is arranged with respect to the weak interaction eigenstate basis: $(\widetilde{B}, \widetilde{W}_3, \widetilde{H}_1, \widetilde{H}_2)$. The angle β is defined in Eq. (7). Three new parameters appear: a supersymmetric Higgs mass term μ and two gaugino Majorana mass terms M_1, M_2 corresponding to masses for $\widetilde{B}$ and $\widetilde{W}_3$. The terms M_1 and M_2 are actually soft-supersymmetry breaking terms which depend on the precise mechanism for supersymmetry breaking. In many models, they are related via:

$$\frac{M_1}{M_2} = \frac{5}{3} \frac{g'^2}{g^2} \tag{11}$$

where the coupling constants on the right-hand side of Eq. (11) are to be taken as running couplings evaluated at a scale $Q^2 \approx O(m_Z^2)$. (For example, Eq. (11) appears naturally in grand unified models. At the grand unified scale, $\frac{5}{3} g'^2 = g^2$ so that at this scale $M_1 = M_2$ which is equal to the common grand unified gaugino mass.)

It is convenient to discuss a simple model in which Eq. (10) can be diagonalized explicitly. Here, we shall consider the model where

$$M_1 = M_2 \equiv M_{\widetilde{\gamma}} \tag{12}$$

$$\mu = 0 \; . \tag{13}$$

Actually, using Eq. (11), we see that $M_1 \neq M_2$ at the electroweak scale. However, as long as $M_1, M_2 \ll m_Z$, it is not a bad approximation to take M_1 and M_2 equal. The advantage of studying the limiting case displayed above is that we may immediately write down the physical neutralino eigenstates:[16]

$$\widetilde{\gamma} = \widetilde{W}_3 \sin\theta_W + \widetilde{B}\cos\theta_W \tag{14}$$

$$\widetilde{H}^0 = \widetilde{H}_1 \sin\beta + \widetilde{H}_2 \cos\beta \tag{15}$$

$$\widetilde{Z}_+ = (\widetilde{W}_3 \cos\theta_W - \widetilde{B}\sin\theta_W)\cos\phi_0 + (\widetilde{H}_1\cos\beta - \widetilde{H}_2 \sin\beta)\sin\phi_0 \tag{16}$$

$$\widetilde{Z}_- = \gamma_5 \left[-(\widetilde{W}_3 \cos\theta_W - \widetilde{B}\sin\theta_W)\sin\phi_0 + (\widetilde{H}_1\cos\beta - \widetilde{H}_2 \sin\beta)\cos\phi_0\right] \tag{17}$$

The corresponding masses are as follows: $M_{\widetilde{\gamma}}$ is the photino mass (hence, the definition given in Eq. (12)); and:

$$M_{\widetilde{H}^0} = 0 \tag{18}$$

$$M_{\widetilde{Z}_\pm} = \left(m_Z^2 + \frac{1}{4} M_{\widetilde{\gamma}}^2 \right)^{1/2} \pm \frac{1}{2} M_{\widetilde{\gamma}} \, . \tag{19}$$

The mixing angle ϕ_0 in Eqs. (16), (17) is given by

$$\tan 2\phi_0 = \frac{2m_Z}{M_{\widetilde{\gamma}}} \, . \tag{20}$$

Thus, in this model $\widetilde{H}^0$ is a pure Higgsino state which is the LSP. The $\widetilde{Z}_\pm$ are mixtures of a pure zino $\widetilde{Z} \equiv \widetilde{W}_3 \cos\theta_W - \widetilde{B} \sin\theta_W$ and Higgsino states; their masses satisfy $M_{\widetilde{Z}_-} \le m_Z$ and $M_{\widetilde{Z}_+} \ge m_Z$. The Dirac matrix γ_5 has been inserted in Eq. (17) so that the mass of the $\widetilde{Z}_-$ is positive as shown in Eq. (19). (This is necessary since the corresponding mass eigenvalue as computed from Eq. (10) is negative. For further details, see Appendix A of Ref. 12.) Typically, one takes $M_{\widetilde{\gamma}}$ small so that in this model the neutralino spectrum satisfies:

$$M_{\widetilde{H}^0} < M_{\widetilde{\gamma}} < M_{\widetilde{Z}_-} < M_{\widetilde{Z}_+} \tag{21}$$

with three of the four neutralinos less massive than the Z^0.

To guide our thinking, we will sometimes further restrict the model by taking $M_{\widetilde{\gamma}} \ll m_Z$ and $v_1 \approx v_2$. (The near equality of the two Higgs vacuum expectation values is common to many low energy supergravity models.[17]) Then, Eqs. (7)–(9), (19) and (20) imply that $-\alpha = \beta = \phi_0 = 45°$ and $m_{H_1^0} = M_{\widetilde{Z}_\pm} = m_Z$. This corresponds to the supersymmetric limit where $(H_1^0, \widetilde{Z}_-, \widetilde{Z}_+, Z^0)$ combine to form one common massive supermultiplet.

3. Phenomenological Consequences of a Light Higgsino

We will study some of the implications of a model whose neutralino stat are given by Eqs. (14)-(17) with mass spectrum as shown in Eq. ... ... let us list in Fig. 1 some important Feynman rules[12,18,19] involving the $\widetilde{H}^0$ (which in the model considered here is the massless state exhibited in Eq. (15)). In diagrams (a) and (b), the factor of $1-\gamma_5$ $(1+\gamma_5)$ corresponds to the production of $\tilde{q}_L(\tilde{q}_R)$. Similar vertices also exist for up and down type leptons. The rules involving the $\widetilde{Z}_-$ depend on the fact that a γ_5 appears in the definition of the state (see Eq. (17)). As a result, a factor of γ_5 appears in diagram (e) involving the scalar H_2^0, but does *not* appear in diagram (f) involving the pseudoscalar H_3^0. A derivation of the Z^0 couplings can be found in the Appendix.

An important feature of these rules is that the $\widetilde{H}^0 q\tilde{q}$ vertex is suppressed by a factor of m_q/m_W. This is a well-known property of the coupling of Higgs bosons to quark pairs ($H^0 q\bar{q}$ vertex), and it occurs here by virtue of the supersymmetry. Note also the appearance of mixing angles in the rules: ϕ_0 (Eq. (20)), β (Eq. (7)) and α (Eq. (9)). In some limiting cases, these angles take on values which may cause certain interactions to vanish. For example, in the supersymmetric limit discussed at the end of Section 2, we saw that $-\alpha = \phi_0 = \beta = 45°$. This implies that in this particular limit, the Feynman rules for the vertices shown in Fig. 1(d), (e) and (f) vanish exactly!

We may now discuss the phenomenological implications of the scenario described above. First, since the photino is no longer the LSP, it is unstable. There are two possible types of decays:

$$(i)\ \tilde{\gamma} \to f\bar{f}\widetilde{H}^0 \quad (f = \text{any fermion such that } m_f < M_{\tilde{\gamma}}/2) \tag{22}$$

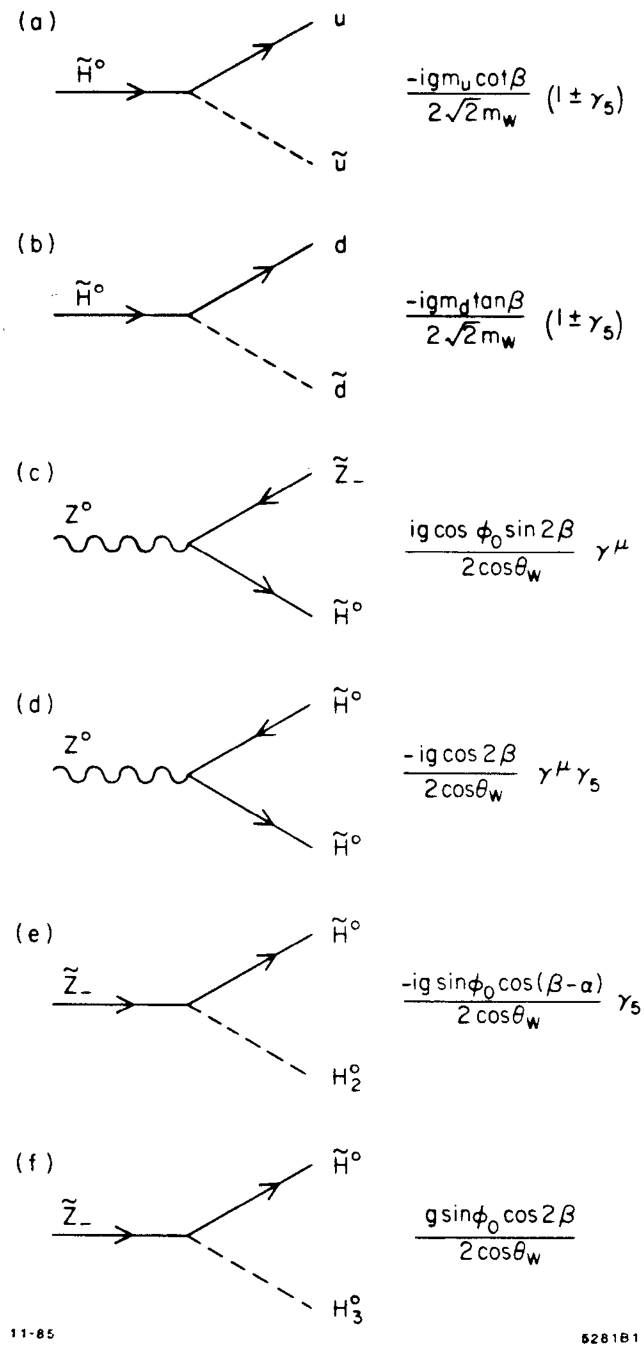

Fig. 1. A partial list of Feynman rules involving the Higgsino $\widetilde{H}^0$, in the model described in the text.

$$(ii)\ \tilde{\gamma} \to \gamma + \widetilde{H}^0 \ . \tag{23}$$

The decay mechanisms are shown in Fig. 2. One can easily show that the process $\tilde{\gamma} \to \gamma + \widetilde{H}^0$ is the dominant mechanism[20,21,18] (unless the photino is very heavy). The reason is that Fig. 2(a) is suppressed due to the smallness of the $\widetilde{H}^0 f \tilde{f}$ coupling (which is proportional to m_f/m_W). On the other hand, in Fig. 1(b), the loop is dominated by the heaviest quark. In fact, in the limit of large m_t, $M_{\tilde{t}}$ with $m_t/M_{\tilde{t}} \sim O(1)$, the rate obtained from Fig. 1(b) approaches a constant non-zero value. The result of a complete calculation (where for definiteness, we take $m_t = M_{\tilde{t}}$ and $\beta = 45°$) is:

$$\tau_{\tilde{\gamma}} \approx 10^{-11} \left(\frac{1\ GeV}{M_{\tilde{\gamma}}} \right)^3 \ sec. \tag{24}$$

which indicates that the photino decay is prompt unless $M_{\tilde{\gamma}}$ is sufficiently light.

How do previous analyses on mass limits of supersymmetric particles change under this scenario? We shall assume that the photino decay is prompt (i.e. $M_{\tilde{\gamma}} \gtrsim O(1 \text{ GeV})$). Consider the process $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$, followed by $\tilde{\gamma} \to \gamma \widetilde{H}^0$. Events of this type will be observed as $e^+e^- \to \gamma\gamma +$ missing energy. Such events have been searched for at PETRA; no signal above background has been seen.[22] This implies that the cross-section for $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$ cannot be too large. Since this process occurs via the exchange of a scalar-electron ($\tilde{e}$), the absence of this process implies a limit on the scalar-electron mass. In Ref. 22, the limit obtained was roughly $M_{\tilde{e}} \gtrsim 100$ GeV. (Note that this limit only applies in the case of an unstable photino which decays before leaving the detector.)

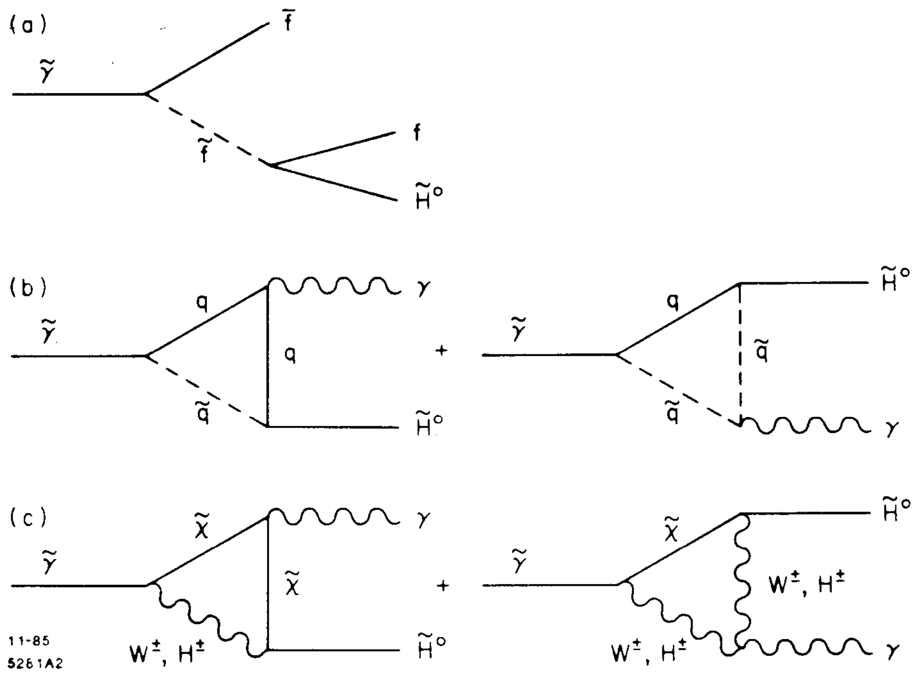

Fig. 2. Mechanisms for Photino Decay. The one-loop process $\tilde{\gamma} \rightarrow \gamma + \widetilde{H}^0$ dominates $\tilde{\gamma} \rightarrow f\bar{f}\widetilde{H}^0$ (where f is any fermion which is kinematically allowed in the decay).

Next, consider the production of scalar-quarks and gluinos at hadron colliders.[23] In previous works, it was assumed that these particle decayed via $\tilde{q} \to q\tilde{\gamma}$ (or $q\tilde{g}$ if kinematically allowed) and $\tilde{g} \to q\bar{q}\tilde{\gamma}$. If the Higgsino is the LSP, one must consider these same decays with $\tilde{\gamma}$ replaced by $\widetilde{H}^0$. But in each case, this involves a $q\tilde{q}\widetilde{H}^0$ vertex which is suppressed by a factor of m_q/m_W. Hence the decays of $\tilde{q}$ and $\tilde{g}$ into photinos will dominate over decays involving Higgsinos. However, the final state photinos will decay via $\tilde{\gamma} \to \gamma\widetilde{H}^0$ with lifetime given by Eq. (24). Thus, *e.g.* for the gluino, the decay chain will be $\tilde{g} \to q\bar{q}\gamma\widetilde{H}^0$, where the $\widetilde{H}^0$ will constitute the missing transverse energy of the event. Two conclusions follow. First, in events with scalar-quark and/or gluino production, less missing energy will result in this scenario as compared with the case where the $\tilde{\gamma}$ is the LSP. This implies that fewer such events would satisfy the cuts and triggers applied by the UA1 collaboration in their search for monojets.[24] Second, such events would contain isolated photons. It is not known at present to what extent such events can be ruled out in the UA1 detector. The end result is that the bounds on scalar-quark and gluino masses are less severe in the case where the $\widetilde{H}^0$ is the LSP. (For a quantitative discussion, see Ref. 23).

The phenomenology of supersymmetric decays of the Z^0 must also be reconsidered. Recalling that $M_{\widetilde{Z}_-} \le m_Z$, we compute:[18]

$$\frac{\Gamma(Z^0 \to \widetilde{Z}_- \widetilde{H}^0)}{\Gamma(Z^0 \to \nu\bar{\nu})} = 2\cos^2\phi_0 \sin^2 2\beta \left(1 - \frac{M^2_{\widetilde{Z}_-}}{m_Z^2}\right)^2 \left(1 + \frac{M^2_{\widetilde{Z}_-}}{2m_Z^2}\right) . \tag{25}$$

If we use Eq. (19) for $M_{\widetilde{Z}_-}$, Eq. (20) for ϕ_0 and take $M_{\tilde{\gamma}} \ll m_Z$, we obtain

$$\frac{\Gamma(Z^0 \to \widetilde{Z}_- \widetilde{H}^0)}{\Gamma(Z^0 \to \nu\bar{\nu})} \approx \frac{3M^2_{\tilde{\gamma}} \sin^2 2\beta}{2m_Z^2} . \tag{26}$$

Typically $\beta \approx 45°$ so that $\sin 2\beta \approx 1$. Thus, if $M_{\widetilde{\gamma}}$ is not too small, the branching ratio for $Z^0 \to \widetilde{Z}_- \widetilde{H}^0$ could be non-negligible. The signature of $Z^0 \to \widetilde{Z}_- \widetilde{H}^0$ could be spectacular at SLC or LEP if the mass of the $\widetilde{Z}_-$ is sufficiently light. The $\widetilde{H}^0$ will not be detected, while the $\widetilde{Z}_-$ will decay into leptons or hadronic jets, resulting in a "one-sided event" (i.e. nearly *all* decay products approximately confined to one-hemisphere of the detector). For a heavier $\widetilde{Z}_-$, the signature will be less striking but should be observable on a statistical basis due to the missing energy resulting from the escaping Higgsinos. In this regard, it is useful to examine in more detail the possible decay channels of the $\widetilde{Z}_-$. In Fig. 3, we exhibit three possible classes of decays:

$$(i) \quad \widetilde{Z}_- \to \widetilde{H}^0 q\bar{q},\ \widetilde{H}^0 \ell^+ \ell^- \tag{27}$$

$$(ii) \quad \widetilde{Z}_- \to \begin{cases} \widetilde{\gamma} q\bar{q},\ \widetilde{\gamma} \ell^+ \ell^- \\ \tilde{g} q\bar{q} \end{cases} \tag{28}$$

$$(iii) \quad \widetilde{Z}_- \to \widetilde{H}^0 H^0 \ . \tag{29}$$

[A fourth channel: $\widetilde{Z}_- \to \widetilde{\chi}^{\pm} q\bar{q}'$, $\widetilde{\chi}^{\pm} \ell\nu$ via $W^{\pm}$ exchange is possible if there exists a chargino state, $\widetilde{\chi}^{\pm}$ lighter than the $\widetilde{Z}_-$.] We shall assume here that $M_{\tilde{\nu}} > M_{\widetilde{Z}_-}$; otherwise, $\widetilde{Z}_- \to \nu\bar{\tilde{\nu}}$ would be the dominant decay channel. If $M_{\tilde{g}} \ll M_{\widetilde{Z}_-}$, then the decay $\widetilde{Z}_- \to \tilde{g} q\bar{q}$ will certainly dominate since it involves a strong interaction $\tilde{g}\tilde{q}q$ vertex. Otherwise, the major decay channels are somewhat model dependent. It is amusing to observe that the sequence $Z^0 \to \widetilde{Z}_- \widetilde{H}^0$ followed by $\widetilde{Z}_- \to \widetilde{H}^0 H^0$, may provide a new mechanism for producing and detecting Higgs bosons (H^0) at SLC and LEP.

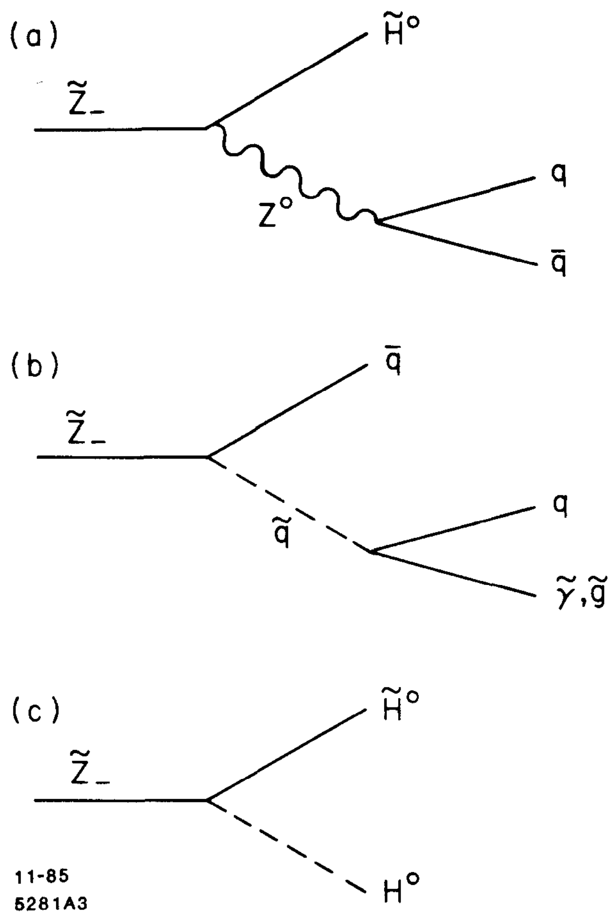

Fig. 3. Mechanisms for $\widetilde{Z}_-$ decay. We do not show the possibility of $\widetilde{Z}_- \rightarrow \widetilde{\chi}^{\pm} q\bar{q}$, which is similar to mechanism (a) via W-exchange (where $\widetilde{\chi}^{\pm}$ is a chargino state). In all diagrams above (except (b) with a final state gluino), we may replace quarks with leptons. We assume that $M_{\tilde{\nu}} > M_{\widetilde{Z}_-}$; otherwise, $\widetilde{Z}_- \rightarrow \nu\bar{\tilde{\nu}}$ would be the dominant decay channel.

We have already remarked that in the minimal supersymmetric model, there necessarily exists a neutral scalar Higgs boson (denoted by H_2^0) with $m_{H_2^0} \leq m_Z$. In fact, from Eq. (8), we see that in the limit of $v_1 = v_2$ (i.e. $\beta = 45°$), $m_{H_2^0} = 0$. This is a tree level result only. If we include radiative corrections *à la* Coleman and E. Weinberg,[25] we find $m_{H_2^0} \simeq O(10 \text{ GeV})$. This is good news for the traditional Higgs boson searches at SLC and LEP. In particular, one would expect to discover such a Higgs boson in a number of ways:[26]

$$(i) \quad Z^0 \to H_2^0 \mu^+ \mu^- \tag{30}$$

$$(ii) \quad Z^0 \to H_2^0 \gamma \tag{31}$$

$$(iii) \quad e^+ e^- \to Z^0 H_2^0 \tag{32}$$

$$(iv) \quad {}^3S_1(t\bar{t}) \to H_2^0 \gamma \tag{33}$$

where ${}^3S_1(t\bar{t})$ is the toponium spin 1 bound state analogous to the ψ. Furthermore, if $\beta = -\alpha = 45°$, then the $\widetilde{Z}_- \widetilde{H}^0 H_2^0$ vertex vanishes and there is no new supersymmetric process. However, it is quite easy to design a model where the $\widetilde{Z}_- \widetilde{H}^0 H_2^0$ vertex is not particularly suppressed.[27,19] For certain values of the H_2^0 mass, the process:

$$Z^0 \to \widetilde{Z}_- \widetilde{H}^0 \to \widetilde{H}^0 \widetilde{H}^0 H_2^0 \tag{34}$$

could well be the dominant mode for Higgs production in Z^0 decays. The dominant decay of H_2^0 would be into $b\bar{b}$ pairs so that the signature of such a process would be a $b\bar{b}$ event (which might be distinguished with a vertex detector) along with substantial missing energy.

In the above discussion, the Higgs boson H_2^0 is a scalar particle and is necessarily lighter than the Z^0. There also exists a pseudoscalar, H_3^0 in these models whose mass is more model dependent—it may be either lighter or heavier than the Z^0. If there exists a light pseudoscalar H_3^0, the above analysis may be repeated with H_3^0 replacing H_2^0. Note in particular that for a pseudoscalar Higgs boson, processes (30) and (32) do not exist because there is no tree level $Z^0Z^0H_3^0$ vertex. (The proof of this claim is easy.[28] Since H_3^0 is a CP-odd state, the only possible gauge-invariant form for the $Z^0Z^0H_3^0$ interaction is $\epsilon_{\mu\nu\alpha\beta}F^{\mu\nu}F^{\alpha\beta}H_3^0$ which is dimension 5 and thus cannot arise in tree level.) Thus, among the processes considered above, only processes (31), (33) and (34) can be used for detecting a light pseudoscalar at a high energy e^+e^- collider.[29] Finally, there exists one new possibility to consider: the decay $Z^0 \to H_2^0H_3^0$. Using the rules given in Ref. 12, I find:

$$\frac{\Gamma(Z^0 \to H_2^0H_3^0)}{\Gamma(Z^0 \to \nu\bar{\nu})} = \frac{1}{2}\cos^2(\beta-\alpha)\, B^3 \tag{35}$$

where B^3 is the usual p-wave suppression factor which occurs for Z^0 decay into a pair of scalars. If $v_1 = v_2$, then $-\alpha = \beta = 45°$ and the rate for $Z^0 \to H_2^0H_3^0$ vanishes. Otherwise, the branching ratio can be non-negligible, in which case this process may provide a good signature for discovering a pair of light Higgs bosons.

Let us return to examine further consequences of a light Higgsino. From Fig. 1(d), the existence of a $Z^0\widetilde{H}^0\widetilde{H}^0$ vertex suggests searching for the process:

$$e^+e^- \to \gamma\widetilde{H}^0\widetilde{H}^0 \tag{36}$$

and using the photon to tag the event (since the Higgsinos will escape undetected). This process is shown in Fig. 4; it is exactly analogous to the classic neutrino counting experiment which has already been used by the ASP Collaboration to set limits for the $\tilde{e}$ and $\tilde{\gamma}$ masses.[30] It is easy to show that:

$$\frac{\Gamma(Z^0 \to \widetilde{H}^0 \widetilde{H}^0)}{\Gamma(Z^0 \to \nu\bar{\nu})} = \cos^2 2\beta \tag{37}$$

which indicates that an $\widetilde{H}^0$ will at best count as one extra neutrino.

We have for the most part neglected the "chargino" states, $\widetilde{\chi}_1^{\pm}$ and $\widetilde{\chi}_2^{\pm}$, (i.e. the mass eigenstates which result from the charged gaugino-Higgsino sector) in the discussions above. One of the charginos may be lighter than the $W^{\pm}$ (the appropriate formulas are similar to Eqs. (19), (20)). For example, there is a $W^-\widetilde{\chi}^+\widetilde{H}^0$ vertex analogous to Fig. 1(c), which would allow for the decay $W^{\pm} \to \widetilde{\chi}^{\pm} + \widetilde{H}^0$. The analysis is parallel to much of the discussion above.

4. Conclusions and Discussion

Even if supersymmetry exists at the electroweak scale, there is at present no strong reason to assume that the photino is the lightest supersymmetric particle. In this paper, a simple model is constructed where the lightest supersymmetric particle is a pure Higgsino state. Because the $\widetilde{H}^0 q\tilde{q}$ vertex is suppressed by a factor m_q/m_W, the phenomenology of a light Higgsino is quite different from that of a light photino. Some of the changes to the "standard" supersymmetric phenomenology are discussed. Some emphasis is placed on signatures of supersymmetry in Z^0 decays. In particular, a new mechanism for the production of a light Higgs boson via $Z^0 \to \widetilde{H}^0\widetilde{H}^0 H^0$ is examined.

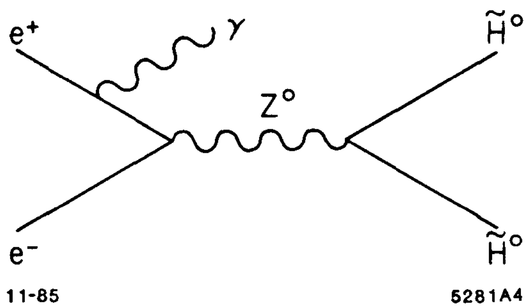

Fig. 4. Detection of Higgsinos using a "neutrino-counting" type experiment. The Higgsinos escape the detector and only the radiative photon is observed.

I shall end with two words of caution. First, certain vertices in the model depend on combinations of mixing angles which can vanish in certain limiting cases. Second, when more realistic models are developed, it could turn out that the lightest supersymmetric particle $\tilde{\chi}^0$ is dominantly made up of a Higgsino component but with some admixture of a gaugino component. The strength of the $\tilde{\chi}^0 q\tilde{q}$ vertex will then crucially depend on the precise amount of this gaugino component. This will have an important bearing on the phenomenology of the supersymmetric model.

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APPENDIX: Z^0 Couplings to Neutralinos: A Tutorial

In this appendix, we demonstrate how to compute the coupling of the Z^0 to a pair of neutralinos.

Consider the general case of a four component fermion:

$$\psi = \begin{pmatrix} \xi_L \\ \eta_R \end{pmatrix} \tag{A1}$$

where ξ_L and η_R may have different $SU(2) \times U(1)$ quantum numbers. If we denote their quantum numbers by T_{3L}, Q and T_{3R}, Q respectively (where the electric charge Q of ξ_L and η_R must be the same), then the Feynman rule for the $\psi\bar{\psi}Z^0$ vertex, denoted by $V(Z^0\bar{\psi}\psi)$, is:

$$\frac{-ig}{2\cos\theta_W}\,\gamma^\mu\;\left[(T_{3L} - Q\sin^2\theta_W)(1-\gamma_5) + (T_{3R} - Q\sin^2\theta_W)(1+\gamma_5)\right]\;. \tag{A2}$$

We shall rewrite this rule in the following form:

$$\frac{-ig}{2\cos\theta_W}\,\gamma^\mu(g_V + g_A\gamma_5) \tag{A3}$$

where

$$g_V = T_{3L} + T_{3R} - 2Q\sin^2\theta_W \tag{A4}$$

$$g_A = T_{3R} - T_{3L}\;. \tag{A5}$$

As an example, for the electron, $T_{3L} = -\frac{1}{2}, T_{3R} = 0$ and $Q = -1$, leading to $g_V = -\frac{1}{2} + 2\sin^2\theta_W$ and $g_A = \frac{1}{2}$ as expected.

It is often convenient to list the fermion content of the theory in terms of left-handed fields only. For example, given two left-handed fermions ξ_L and η_L of opposite charge, I can construct the Dirac fermion of Eq. (1) by identifying $\eta_R = (\eta_L)^c$, i.e. the charge-conjugated state. I would then identify $T_{3L} = T_3(\xi_L)$ and $T_{3R} = T_3(\eta_R) = -T_3(\eta_L)$. This discussion is also valid for a Majorana fermion, i.e. where $\xi = \eta$ in Eq. (A1). It then follows that for a Majorana fermion, $T_{3R} = -T_{3L}$ and $Q = 0$.

We illustrate these ideas by computing the Feynman rule for the $Z^0\widetilde{H}^0\widetilde{H}^0$ vertex. Using Eq. (15), this vertex can be expressed as,

$$V(Z^0\widetilde{H}^0\widetilde{H}^0) = \sin^2\beta\, V(Z^0\widetilde{H}_1\widetilde{H}_1) + \cos^2\beta\, V(Z^0\widetilde{H}_2\widetilde{H}_2) \qquad (A6)$$

where $\widetilde{H}_1$ and $\widetilde{H}_2$ are the (left-handed) weak interaction eigenstates (with definite $SU(2)\times U(1)$ quantum numbers). Note that $V(Z^0\widetilde{H}_1\widetilde{H}_2) = 0$ since the Z^0 does not couple off-diagonally to weak interaction eigenstates at tree level. The states $\widetilde{H}_1$ and $\widetilde{H}_2$ are neutral Majorana fermions; therefore $T_{3R} = -T_{3L}$ in each case. In supersymmetry, $\widetilde{H}_1$ and $\widetilde{H}_2$ have opposite hypercharge, and $T_{3L}(\widetilde{H}_1) = \frac{1}{2}$ and $T_{3L}(\widetilde{H}_2) = -\frac{1}{2}$. Using the above results, it then follows that $g_V = 0$ and $g_A = \cos^2\beta - \sin^2\beta$ which reproduces the rule stated in Fig. 1(d).

The $Z^0\widetilde{H}^0\widetilde{Z}_-$ vertex can be obtained in a similar manner. Here there are a few subtleties. We must keep track of which fermion is outgoing in the vertex (that fermion will have a *bar* over it), and we must pay attention to the factor of γ_5 in Eq. (17). It then follows from Eq. (17) that:

$$V(Z^0\overline{\widetilde{H}}{}^0\widetilde{Z}_-) = \cos\phi_0\sin\beta\cos\beta\left[V(Z^0\overline{\widetilde{H}}_1\gamma_5\widetilde{H}_1) - V(Z^0\overline{\widetilde{H}}_2\gamma_5\widetilde{H}_2)\right] \ . \qquad (A7)$$

The presence of the factor of γ_5 in Eq. (A7) changes the Feynman rule given in

Eq. (A3) to:

$$\frac{-ig}{2\cos\theta_W}\,\gamma^\mu(g_V\gamma_5 + g_A)\ . \tag{A8}$$

Otherwise, the calculation is analogous to the one above, and we find $g_V = 0$ and $g_A = -2\cos\phi_0\sin\beta\cos\beta$. Inserting this into Eq. (A8), we obtain the rule given in Fig. 1(c).

The method described above is of course applicable to the Z^0 coupling to any fermion pair. As one last example, consider the charged left-handed Higgsino states $\widetilde{H}_1^-$ and $\widetilde{H}_2^+$ with $T_{3L} = -\frac{1}{2}$ and $+\frac{1}{2}$ respectively. A Dirac fermion $\widetilde{\chi}^+$ can be constructed:

$$\widetilde{\chi}^+ = \begin{pmatrix} \widetilde{H}_2^+ \\ (\widetilde{H}_1^-)^c \end{pmatrix}\ . \tag{A9}$$

In this example, $T_{3L} = T_{3R} = \frac{1}{2}$ so that the $Z^0\widetilde{\chi}^+\widetilde{\chi}^-$ coupling is purely vector. From Eq. (A4), it follows that $g_V = 1 - 2\sin^2\theta_W$.

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(14)). Nevertheless, as long as $M_1, M_2 \ll m_Z$, the deviation from the simple model discussed in the text is small, and all results which follow remain approximately correct.

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28. I can give a second proof which is a little more subtle. If I omit the fermions from the Glashow-Weinberg-Salam model (with two complex Higgs doublets and a CP–invariant potential), then the theory conserves P and C separately. A consistent set of quantum numbers can be found for the particles in the model; *e.g.*, the Z^0 has $J^{PC} = 1^{--}$, H_1^0 and H_2^0 are 0^{++} and H_3^0 is 0^{+-}. The $Z^0Z^0H_3^0$ vertex is C–violating and therefore does not exist! (If I reintroduce the fermions as usual, then the Z^0 becomes a mixture of 1^{--} and 1^{++}, and the H_3^0 becomes a mixture of 0^{+-} and 0^{-+}. Only the 0^{-+} can contribute to the $H_3^0 f\bar{f}$ coupling, thereby allowing a $Z^0Z^0H_3^0$ vertex via the triangle diagram with fermions in the loop.)
29. It may be quite difficult to detect a light pseudoscalar Higgs boson at SLC or LEP. The branching ratio for $Z^0 \to H^0\gamma$ is $O(10^{-6})$ for a scalar Higgs boson and even smaller for a pseudoscalar Higgs boson (where only the quark loop contributes in the one-loop Feynman diagram). Hence, process (31) is not a realistic way to observe the Higgs boson.
30. D. L. Burke, contribution to these proceedings.