PHYSICAL REVIEW D VOLUME 37, NUMBER 3 1 FEBRUARY 1988

Supersymmetry phenomenology and the nature of the lightest supersymmetric particle

Manuel Drees, C. S. Kim, and Xerxes Tata
Physics Department, University of Wisconsin, Madison, Wisconsin 53706
(Received 24 July 1987)

We show that in addition to the photinolike and Higgsino-like states usually considered as candidates for the lightest supersymmetric particle (LSP), there is yet one more state in the neutralino sector which could be the LSP in minimal supergravity models. In this paper, we study how the phenomenology of supersymmetry is altered depending on the nature of the LSP. Within the framework of minimal supergravity models, we demonstrate that if the mass of the LSP is small compared to that of the W boson, the lightest chargino and the next to lightest neutralino states are always lighter than M_W and M_Z independent of the soft-supersymmetry-breaking gaugino masses. We show that the W- and Z-boson decay widths into these particles depends quite significantly on the nature of the LSP. We further show that the bound on the scalar-electron mass of the ASP detector at the SLAC storage ring PEP is considerably weakened if the LSP is the new state discussed above, even if this state is dominantly a gaugino. Finally, we briefly study the phenomenology of the LSP in the context of superstring-inspired E(6) models.

I. INTRODUCTION

In the last five years, a considerable amount of effort[1] has been devoted to the construction and phenomenological implications of $N=1$ supergravity models. These efforts have gained renewed interest with the discovery[2] of anomaly-free superstring theories, which are a candidate for a unified description of gravitational and Yang-Mills interactions. It has been shown[3] that their compactifications to four dimensions can lead to an $N=1$ supergravity theory in the low-energy limit.

One feature of most supergravity models, including superstring-inspired ones, is the existence of a conserved discrete symmetry called R parity; the normal fermions and Higgs and gauge bosons are R even, while their superpartners are R odd. It then follows that the lightest supersymmetric particle (LSP) is stable; all other superpartners decay into the LSP plus one or several R-even particles. Therefore, the experimental signals for the production of any R-odd particle strongly depends on the nature of the LSP.

It is known that the LSP must be electrically and color neutral;[4] otherwise an abundance of exotic nuclei would have been observed.[5] In presuperstring models,[1,6] which are essentially a direct supersymmetrization of the standard model, there are thus two obvious candidates for the LSP: The sneutrino $\tilde{\nu}$ and the lightest eigenstate $\tilde{Z}_1$ of the neutral gauge–Higgs-fermion sector.

The first possibility, which has been discussed in Refs. 7, now seems somewhat disfavored, since there exist quite restrictive lower bounds on masses of charged sleptons[8] and squarks;[9] in most models the sleptons have similar masses, so that a light sneutrino seems unlikely. Although superstring-inspired E(6) models in principle allow for the required mass splitting, a substantial fine-tuning is needed to create a light sneutrino, since in these models scalars tend to be heavy.[10] We will therefore not pursue the possibility that a sneutrino is the LSP any further. In minimal supergravity models[1,6] this leaves us with the $\tilde{Z}_1$.

There are three distinct scenarios which lead to a light $\tilde{Z}_1$. In the first and most widely studied[1,11] scenario it is assumed that the soft-supersymmetry-breaking gaugino masses are small compared to the mass of the Z boson, M_Z. In this case $\tilde{Z}_1$ is approximately the superpartner of the photon, the photino. In the second scenario[12] one postulates a small supersymmetric Higgs-boson mass, in which case $\tilde{Z}_1$ is dominantly a mixture of the two Higgsino current states. Finally the $\tilde{Z}_1$ can be made light by an explicit cancellation among the terms that contribute to the determinant of the neutral-gaugino–Higgsino mass matrix. At first glance this looks like an unnatural fine-tuning; it turns out, however, that there is a rather large region of parameter space for which such a light $\tilde{Z}_1$ results. Furthermore, unlike the first two scenarios this last scenario allows all entries of the neutralino mass matrix to be of the same order of magnitude. Moreover, much of this region of parameter space is allowed by current experimental data, and if the LSP is indeed such a light $\tilde{Z}_1$, conventional supersymmetry (SUSY) phenomenology is quite altered. We have analyzed the phenomenology in each of these cases, with a particular emphasis on the third scenario where the terms in the determinant of the neutralino mass matrix approximately cancel, since this region has not been well studied in the literature.

In superstring-inspired models there exist six additional R-odd neutral fermions[13] which usually do not mix[14] with the gaugino-Higgsino sector. It has recently been shown[15] that in models without an intermediate scale the lightest of these particles cannot be heavier than about 115 GeV and is often much lighter. In fact it is possible[13] that the lightest of these exotic fermions is the LSP. In this case, the phenomenology of the LSP, and thus of all R-odd particles, strongly depends on the unknown superpotential couplings of the exotic E(6) fermions. This makes it impossible to make strong quantitative statements; it is, however, possible to qualitatively classify the various possible cases and the emerging signatures. This discussion forms the second part of our paper.

The rest of this paper is organized as follows. In Sec. II we discuss in some detail the space of the parameters that enter the chargino and neutralino mass matrices, as well as existing experimental bounds. In Sec. III we show that if the $\tilde{Z}_1$ is light, the mass of the second-lightest neutralino is always smaller than M_Z, while the lightest chargino is not heavier than M_W, independent of the soft-SUSY-breaking gaugino masses. In Sec. IV we discuss decays of W and Z bosons into neutralinos and charginos, with special emphasis on the region of parameter space that leads to the new light neutralino. In Sec. V we reexamine the existing bound of the ASP detector at the SLAC storage ring PEP on $m_{\tilde{Z}_1}$ as a function of the selectron mass. We show that this bound is substantially stronger for the photino than for the new state (the third scenario of the previous paragraph). In Sec. VI we qualitatively analyze the case that the LSP is a neutral exotic E(6) fermion. Finally, Sec. VII contains a summary of our findings.

II. NEUTRALINOS AND CHARGINOS IN $N=1$ SUPERGRAVITY MODELS

Minimal supergravity models contain four neutral R-odd Majorana fermions: The Higgsinos $\bar{h}^0$ and h^0, which are the superpartners of the $Y=+\frac{1}{2}$ and $Y=-\frac{1}{2}$ Higgs bosons $\bar{H}^0$ and H^0, and the neutral SU(2) and $U(1)_Y$ gauginos λ_3 and λ_0. The neutralino mass matrix in the basis $(\bar{h}^0, h^0, \lambda_3, \lambda_0)$ is given by[1,11,17]

$$M_{(\mathrm{neutral})}=\begin{bmatrix} 0 & -2m_1 & \frac{1}{\sqrt{2}}g\bar{v} & -\frac{1}{\sqrt{2}}g'\bar{v} \\ -2m_1 & 0 & -\frac{1}{\sqrt{2}}gv & \frac{1}{\sqrt{2}}g'v \\ \frac{1}{\sqrt{2}}g\bar{v} & -\frac{1}{\sqrt{2}}gv & \mu_2 & 0 \\ -\frac{1}{\sqrt{2}}g'\bar{v} & \frac{1}{\sqrt{2}}g'v & 0 & \mu_1 \end{bmatrix} . \qquad (2.1)$$

Here $2m_1$ is the supersymmetric Higgs-boson mass, $\bar{v}\equiv\langle\bar{H}^0\rangle$, $v\equiv\langle H^0\rangle$, and μ_2 and μ_1 are soft breaking SU(2) and $U(1)_Y$ gaugino masses. In supergravity grand unified theories (GUT's) with canonical kinetic-energy terms for gauge superfields the latter two parameters are related to the gluino mass $|\mu_3|$ by

$$\mu_1=\frac{5}{3}\tan^2\theta_W\mu_2=\frac{5}{3}\frac{g'^2}{g_S^2}\mu_3 , \qquad (2.2)$$

where g', g, and g_S are the $U(1)_Y$, SU(2), and SU(3) gauge couplings, respectively. We will assume this relation throughout our paper.

In the charged sector, the model contains two R-odd Dirac fermions:[11,17] $\lambda\equiv(\lambda_1-i\lambda_2)/\sqrt{2}$ and $\chi\equiv P_L h - P_R\bar{h}$. The chargino mass matrix contains the same parameters that enter the neutralino mass matrix. In the basis (λ,χ) it reads

$$M_{(\mathrm{charged})}=\begin{bmatrix} \mu_2 & gv \\ g\bar{v} & 2m_1 \end{bmatrix} . \qquad (2.3)$$

Thus all masses and mixing angles in the chargino and neutralino sectors can be parametrized in terms of μ_2, $2m_1$, and $\bar{v}/v$. Following the notation of Refs. 17 and 18 we label the physical neutralino and chargino states by $\tilde{Z}_i$ and $\tilde{W}_{+/-}$, with $\tilde{Z}_1$ and $\tilde{W}_-$ being the lightest states in the corresponding sector. We denote the $\bar{h}^0$, h^0, λ_3, and λ_0 components of $\tilde{Z}_i$ by $v_1^{(i)}$, $v_2^{(i)}$, $v_3^{(i)}$, and $v_4^{(i)}$, respectively.

As discussed in the Introduction, there are three different scenarios that lead to small values of $m_{\tilde{Z}_1}$. This can most easily be seen from the determinant of the neutralino mass matrix, which is given by

$$\det M_{(\mathrm{neutral})}=\frac{2m_1}{3}g'^2\mu_2(8v\bar{v}-2m_1 5\mu_2/g^2) . \qquad (2.4)$$

The first possibility[11] to obtain a massless $\tilde{Z}_1$ is to set $\mu_2=0$; in this case $\tilde{Z}_1$ is the photino, i.e, $\mathbf{v}^{(1)}=(0,0,\sin\theta_W,\cos\theta_W)$. This case has been widely studied in the literature.[1,11] By virtue of Eq. (2.2) a light photino implies a light gluino. In view of the fact that the UA1 Collaboration has been able[9] to exclude gluino masses below 53 GeV the allowed region for this case is rather restricted.

Another way to achieve a small $m_{\tilde{Z}_1}$ is to choose[12] $2m_1$ to be small. In the extreme case, $2m_1=0$, $\tilde{Z}_1$ is a pure Higgsino, i.e., $\mathbf{v}^{(1)}=(\cos\beta,\sin\beta,0,0)$ where $\tan\beta\equiv\bar{v}/v$. This solution is, however, somewhat problematic in minimal models. The reason is that $2m_1$, multiplied by a soft-breaking parameter B, also enters the Higgs potential; one thus expects $2m_1 B$ to be of order M_W^2, which is hard to achieve if $2m_1$ is very small. One solution[12] is to introduce a nonminimal hidden sector which effectively decouples B from the other soft-breaking parameters. It should also be mentioned that, unless $|\tan\beta|=1$, even small nonvanishing values of $2m_1$ suffice to dramatically change[19] the character of $\tilde{Z}_1$ by introducing nonvanishing $v_3^{(1)}$ and $v_4^{(1)}$, which dominate the couplings of $\tilde{Z}_1$ to matter. The existence of a light Higgsino-like state thus seems unlikely. In contrast, a photinolike eigenstate occurs for a substantial range of μ_2 if the signs of the parameters entering the mass matrix (2.1) are chosen such that the two terms in the determinant (2.4) add rather than cancel.

Finally, $m_{\tilde{Z}_1}$ can be set to zero by explicitly canceling the two terms in Eq. (2.4), i.e., by requiring

$$2m_1\mu_2=\tfrac{8}{5}M_W^2\sin(2\beta)\ . \tag{2.5}$$

In this case the $\tilde{Z}_1$ is a complicated mixture of all four current states:

$$\mathbf{v}^{(1)}=\frac{1}{N}\left[\frac{1}{\sqrt{2}},-\frac{1}{\sqrt{2}}\tan\beta,-\frac{M_W}{\mu_2}\sqrt{2}\sin\beta,\frac{M_W}{\mu_2}\sqrt{2}\sin\beta\frac{3}{5\tan\theta_W}\right], \tag{2.6}$$

where N is determined by the requirement $\|\mathbf{v}^{(1)}\|=1$.

We should clarify at this point that there is no compelling reason to expect the LSP to be light (compared to the W boson), let alone massless. This is, however, the only case of phenomenological interest at energies accessible at present or near future colliders such as the Fermilab Tevatron, KEK TRISTAN, Stanford Linear Collider (SLC), and the CERN LEP. We use the three solutions (discussed above) to the equation $m_{\tilde{Z}_1}=0$ as a guide to determining the regions of parameter space that lead to a light LSP. For brevity we shall characterize each of these regions by the particular solution to the equation $m_{\tilde{Z}_1}=0$. Of course, by varying the SUSY model parameters one can continuously go from one region to the other, so that no other significance is attached to these "solutions."

In the minimal model considered here bounds on the neutralino sector cannot only be inferred by experiments that search for these particles, like the ASP experiment,[16] since from Eqs. (2.2) and (2.3) it follows that bounds on chargino and gluino masses constrain the same parameters that determine masses and mixings of the neutralinos. This is illustrated in Figs. 1(a) and 1(b) for $\tan\beta=1$ and 2.5, respectively. In our calculations we used the values $M_W=83$ GeV, $\sin^2\theta_W(M_W)=0.22$, and $\alpha_S(M_W)=0.136$, corresponding to $\Lambda_{QCD}=200$ MeV and 5 quark flavors. Note that we can without loss of generality assume $2m_1$, v, and $\bar{v}$ to be positive;[18] in this case we have to consider both signs for μ_3, with $m_{\tilde{g}}=|\mu_3|$. In these figures the solid curves are lines of $m_{\tilde{Z}_1}=12.5$ GeV, which is the kinematic limit for the ASP experiment.[16] The photinolike solution is characterized by[11] an almost constant $|\mu_3|\simeq 80$ GeV. For $\tan\beta=1$ the Higgsino state is characterized by $2m_1=m_{\tilde{Z}_1}$, whereas the corresponding state for $\omega=2.5$ (which is no longer a Higgsino, as discussed earlier) is characterized by[19] $2m_1\simeq m_{\tilde{Z}_1}/\sin(2\beta)=18$ GeV. Finally the two hyperbolas are contours of $m_{\tilde{Z}_1}=12.5$ GeV, where $\tilde{Z}_1$ is a light state that approximately satisfies Eqs. (2.5) and (2.6);

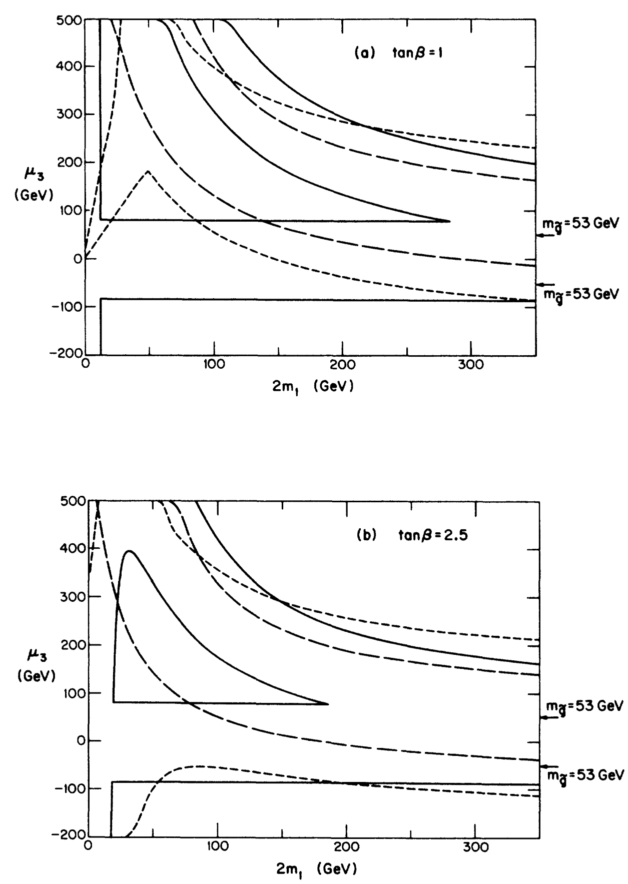

FIG. 1. Various regions in the space of parameters that enter the neutralino and chargino mass matrices for (a) $\tan\beta=1$ and (b) $\tan\beta=2.5$. The solid curves are lines of constant $m_{\tilde{Z}_1}=12.5$ GeV, which is the kinematic limit of the ASP experiment. The region between the two long-dashed curves is excluded since here the PETRA bound $m_{\tilde{W}_-}>22$ GeV is violated. The decay $W\to\tilde{W}_-\tilde{Z}_2$, which leads to spectacular three lepton events at the CERN $Sp\bar{p}S$ and the Tevatron, is kinematically possible in the region between the short-dashed curves. The arrows indicate the UA1 bound $m_{\tilde{g}}\gtrsim 53$ GeV.

there are two solutions since nonzero values of $m_{\tilde{Z}_1}$ can be obtained by either decreasing or increasing μ_2 from the value given by Eq. (2.5), which leads to $m_{\tilde{Z}_1}=0$.

It is interesting to note that two of these three curves sometimes meet, most notably at the corners of the triangle in Fig. 1(b). At the points with $\mu_3 \simeq 80$ GeV, there are two degenerate light neutralinos, while at the upper end there is only one light state; it is equally far away from being a pure Higgsino and the massless state of Eq. (2.6).

The region between the two long-dashed curves is excluded since here the lighter chargino is lighter than 22 GeV, which violates the DESY PETRA bound.[8] Thus, the solution where μ_2 is slightly smaller than given by Eq. (2.5), i.e., the lower of the $m_{\tilde{Z}_1}=12.5$ GeV hyperbolas, is always ruled out. As is well known,[11] situations with a very small μ_2 and large $2m_1$, or vice versa, are also excluded by the bound $m_{\tilde{W}_-} \geq 22$ GeV. Furthermore, many situations with a light photinolike $\tilde{Z}_1$ are excluded by the UA1 bound[9] on the gluino mass which excluded the region $m_{\tilde{g}} \lesssim 53$ GeV.

Additional bounds may be derived[20,21] from the nonobservation of $W \to \tilde{W}_- + \tilde{Z}_2$ and $Z \to \tilde{W}_- \tilde{W}_-$ decays which would lead to spectacular trilepton and acollinear dilepton events. The first decay, which has the largest rate at the CERN $Sp\bar{p}S$, is kinematically allowed in the region between the short-dashed curves in Figs. 1. Although existing experiments are probably only sensitive[20] up to $m_{\tilde{W}_-}+m_{\tilde{Z}_2} \simeq 75$ GeV, these curves indicate that a substantial region with a light photinolike or Higgsino-like $\tilde{Z}_1$ and positive μ_3 may already be ruled out.

In principle, further bounds[20,21] can be derived from the monojet events that originate from the decay $W \to \tilde{W}_- \tilde{Z}_1$, followed by the hadronic decay of $\tilde{W}_-$, in much the same way as the UA1 Collaboration has obtained bounds on the mass of sequential heavy lepton.[22] There are, however, severe standard-model backgrounds to these events so that the exact determination of this bound may be more difficult. For this reason we have not included this bound in the figure. We will discuss gauge-boson decays in more detail in Sec. IV.

Although the results of this section have been derived within the framework of the minimal supergravity model with only four neutralinos they are also valid in the simplest superstring-inspired E(6) models with one additional $\tilde{U}(1)$ gauge group and one SU(2)-singlet Higgs boson N. In these models the supersymmetric Higgsino mass arises from the superpotential coupling $\lambda H\bar{H}N$ after spontaneous breaking of the $\tilde{U}(1)$, i.e., $2m_1=\lambda\langle N\rangle$. The chargino sector is the same as in minimal supergravity. Furthermore, the extra fermions in the neutralino sector decouple from the fermions of minimal supergravity. The reason is that the new neutral gauge boson Z' has to be heavier than the standard Z boson since otherwise its effects would have been seen[23] in existing neutral-current data. Furthermore, in models without an intermediate scale the Z' has to be heavier[24] than 330–400 GeV in order not to destroy successful standard-model predictions for nucleosynthesis. This leads to the effective decoupling[18] of the additional neutralinos from the four neutralinos of the minimal model, leaving us with the mass matrix (2.1), where $2m_1$ has to be replaced by $\lambda\langle N\rangle$. Since the superpotential coupling λ is expected[25,10] to be not much smaller than ~ 0.1 and $M_{Z'} \simeq 5g'\langle N\rangle/3\sqrt{2}$ for $M_{Z'}^2 \gg M_Z^2$, this means that the model predicts a rather large effective value for $2m_1$, which excludes the possibility of a light Higgsino-like state. At least in the region of small $m_{\tilde{Z}_1}$ this is the only difference between minimal and superstring-inspired models. In the next three sections, which discuss phenomenological implications of a light $\tilde{Z}_1$, we will therefore stick to the case of minimal supergravity.

III. BOUNDS ON $m_{\tilde{W}_-}$ and $m_{\tilde{Z}_2}$

It is well known[26] that if the soft-SUSY breaking gaugino masses are small compared with M_W (so that $\tilde{Z}_1 \approx \tilde{\gamma}$ with a small mass), there are light states $\tilde{W}_-$ and $\tilde{Z}_2$ in the chargino and neutralino sectors with $m_{\tilde{W}_-} \leq M_W$ and $m_{\tilde{Z}_2} \leq M_Z$. In this section we study the corresponding limits on $\tilde{Z}_2$ mass when $\tilde{Z}_1$ is either a light Higgsino or the light state described by Eqs. (2.5) and (2.6). Unlike the case of the light photino, which can be analyzed in a model-independent fashion, our considerations are limited to the minimal model, although similar results hold for the E(6) superstring models with a large Z'-boson mass.

In the minimal model the bound on $m_{\tilde{Z}_2}$ for the case $\mu_1=\mu_2=0$ is even stronger than M_Z. One simple method for deriving upper bounds on the absolutely smallest eigenvalue $E_{\min}$ of a given matrix M makes use of the inequality $E_{\min} \leq \|M\mathbf{x}\|$ that holds for all unit vectors $\mathbf{x}$. In the present case M is the 3×3 neutralino mass matrix in the Z-ino–Higgsino basis (obtained by truncating the photino components of the original matrix),[11] which is convenient for the discussion of the $m_{\tilde{\gamma}}=0$ case; it reads

$$M'_{(\text{neutral})} = \begin{pmatrix} 0 & M_Z\cos\beta & -M_Z\sin\beta \\ M_Z\cos\beta & 0 & -2m_1 \\ -M_Z\sin\beta & -2m_1 & 0 \end{pmatrix} . \tag{3.1}$$

Using the vector $\mathbf{x}_1=(\cos\alpha,\sin\alpha,0)$ with $\tan\alpha=-M_Z\sin\beta/2m_1$, one finds $m_{\tilde{Z}_2} \leq M_Z\cos\beta$; similarly one can show that $m_{\tilde{Z}_2} \leq M_Z\sin\beta$ by use of the vector $\mathbf{x}_2=(\cos\alpha,0,\sin\alpha)$ with $\tan\alpha=M_Z\cos\beta/2m_1$.

The bound on the $\tilde{W}_-$ mass can also be obtained from the matrix (2.3). For $\mu_2=0$ it is obvious that the smallest eigenvalue is bounded by the norm of any row or column. We thus have

$$m_{\tilde{W}_-} \leq \sqrt{2}M_W\min(\cos\beta,\sin\beta) ,$$
$$m_{\tilde{Z}_2} \leq M_Z\min(\cos\beta,\sin\beta) \quad (\mu_1=\mu_2=0) . \tag{3.2}$$

The bounds (3.2) have been derived in the limiting situation $m_{\tilde{\gamma}}=0$; the corresponding bounds for finite values of $m_{\tilde{Z}_1}$ are shown in Fig. 2(a) for the case $2m_1\mu_3<0$, since this ensures that the $\tilde{Z}_1$ is photinolike even for $m_{\tilde{Z}_1}$ as large as 40 GeV. It is interesting to note that $m_{\tilde{Z}_2}^{\max}$ rises faster with $m_{\tilde{Z}_1}$ than $m_{\tilde{W}_-}^{\max}$; eventually the two bounds become identical.

Similar bounds can be derived for the case of small $2m_1$. It is well known[1] that the eigenvalues of the chargino mass matrix are invariant under the exchange of $2m_1$ and μ_2; thus the bounds that have been derived for the case $\mu_2=0$, which implies a massless photino, are also valued in the case that the $\tilde{Z}_1$ is a massless Higgsino, i.e., $2m_1=0$.

However, the bound on $m_{\tilde{Z}_2}$ is substantially altered. To see this we first observe that in the case $2m_1=0$ all neutralino masses are independent of β. The bounds on $m_{\tilde{Z}_2}$ can, therefore, be conveniently obtained from the simple case $\tan\beta=1$. The mass matrix for the three massive neutralinos is then given by [in the basis $(1/\sqrt{2})(h_0-\bar{h}^0),\lambda_3,\lambda_0$]

$$M''_{(\text{neutral})}=\begin{bmatrix} 0 & M_W & -M_W\tan\theta_W \\ M_W & \mu_2 & 0 \\ -M_W\tan\theta_W & 0 & \mu_1 \end{bmatrix} . \quad (3.3)$$

By considering the vector $\mathbf{x}_3=(0,\sin\theta_W,\cos\theta_W)$ we find

$$m_{\tilde{Z}_2}^2 \le \mu_2^2\sin^2\theta_W(1+\tfrac{25}{9}\tan^2\theta_W)\equiv a\mu_2^2 . \quad (3.4)$$

On the other hand the vector $\mathbf{x}_4=(\cos\alpha,\sin\alpha,0)$, with $\tan\alpha=-M_W/\mu_2$, gives us the bound

$$m_{\tilde{Z}_2}^2 \le \frac{M_W^2}{\mu_2^2+M_W^2}(M_W^2+\tan^2\theta_W\mu_2^2) . \quad (3.5)$$

Note that the bound (3.4) increases monotonically with μ_2 while the bound (3.5) decreases. $m_{\tilde{Z}_2}$ will therefore be maximal if μ_2 is chosen such that these two bounds coincide. This gives, as the final result,

$$m_{\tilde{Z}_2}^2 \le \frac{M_W^2}{2}(-b+\sqrt{4a+b^2}),\quad b=a-\tan^2\theta_W ,$$
$$m_{\tilde{W}_-} \le \sqrt{2}M_W\min(\cos\beta,\sin\beta)\quad (2m_1=0) . \quad (3.6)$$

Numerically the bound (3.6) on $m_{\tilde{Z}_2}$ is very close to the actual maximal $m_{\tilde{Z}_2}$ for $2m_1=0$, obtained by numerically scanning the parameter space (63 GeV vs 59 GeV for $M_W=83$ GeV, $\sin^2\theta_W=0.22$).

The corresponding bounds for $2m_1\neq0$ are shown in Fig. 2(b). The results for $m_{\tilde{W}_-}^{\max}$ are very similar to those of Fig. 2(a), if one rescales $m_{\tilde{Z}_1}$ [Fig. 2(b)]$=m_{\tilde{Z}_1}$[Fig. 2(a)]$\sqrt{a}/\sin(2\beta)$. This follows from the symmetry of $m_{\tilde{W}_-}$ under the exchange $2m_1\leftrightarrow\mu_2$ and the fact that in Fig. 2(a) one has $m_{\tilde{Z}_1}\simeq\sqrt{a}\,\mu_2$, while the $m_{\tilde{Z}_1}$ of Fig. 2(b)

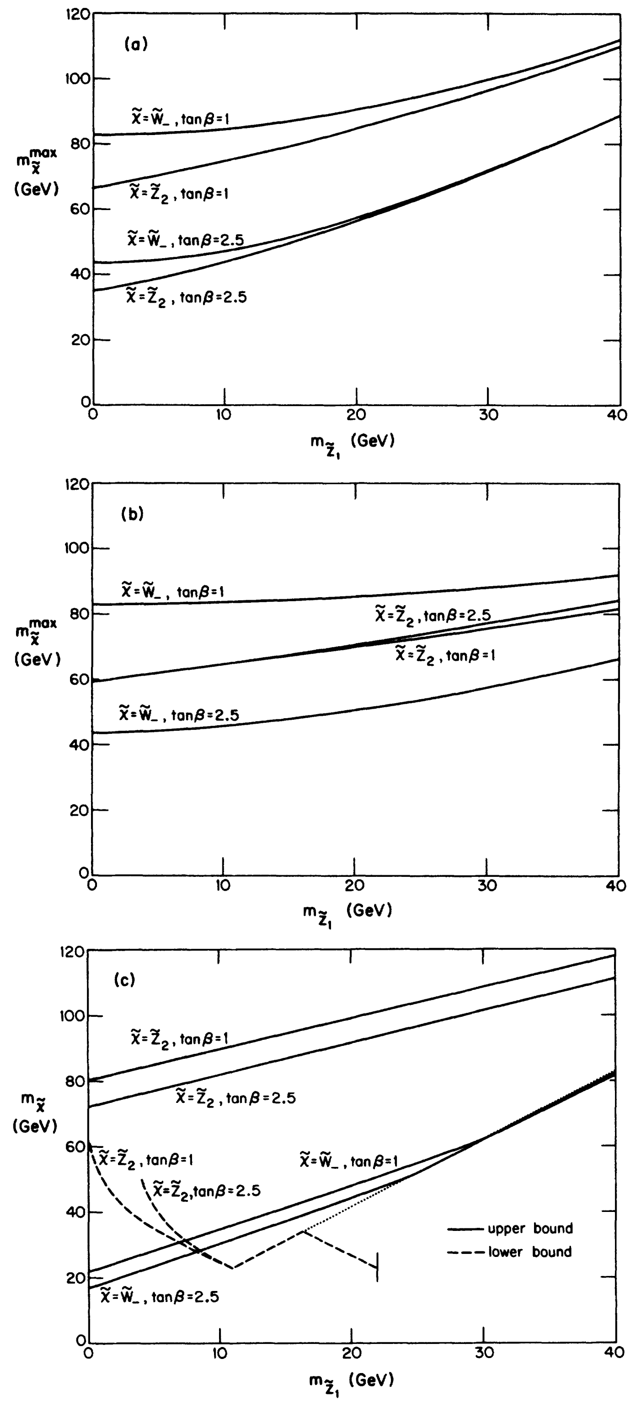

FIG. 2. Upper bounds on $m_{\tilde{W}_-}$ and $m_{\tilde{Z}_2}$ as a function of $m_{\tilde{Z}_1}$ (a) in the region of small and negative μ_2 and positive $2m_1$, (b) in the region of small and negative $2m_1$ and positive μ_2, and (c) in the region $2m_1\ge1.6M_W^2\sin(2\beta)/\mu_2$; see Eq. (2.5). We show results for $\tan\beta=1$ and 2.5, respectively. The dashed curves in (c) show the lower bound on $m_{\tilde{Z}_2}$, which emerges from the condition $m_{\tilde{W}_-}\ge22$ GeV. For $m_{\tilde{Z}_1}>22$ GeV there is no lower bound on $m_{\tilde{Z}_2}$ unless, as discussed in the text, we also bound $m_{\tilde{g}}$ from above. The dotted curve is the lower bound on $m_{\tilde{Z}_2}$ for $m_{\tilde{g}}\lesssim2$ TeV. No corresponding bound exists in cases (a) and (b).

is approximately given by $2m_1\sin(2\beta)$. The plot also shows that $m_{\tilde{Z}_2}^{\max}$ remains almost independent of β as long as $m_{\tilde{Z}_1}$ is not much larger than 40 GeV. In general one sees that in the case of small $2m_1$, $m_{\tilde{Z}_2}^{\max}$ and $m_{\tilde{W}_-}^{\max}$ rise much more slowly with $m_{\tilde{Z}_1}$ than in the case of small μ_1 and μ_2.

Finally, we turn to the case where both $2m_1$ and μ_2 are nonzero, but Eq. (2.5) is satisfied, i.e., $\mu_2=0.8M_W^2\sin(2\beta)/m_1$. By multiplying the chargino mass matrix (2.3) from the right with $\mathbf{x}_5=(\cos\theta,-\sin\theta)$, with

$$\tan\theta=2m_1\sqrt{5}/(4M_W\cos\beta)\ ,$$

one finds

$$m_{\tilde{W}_-}^2\le 2M_W^2(1-\sqrt{8/5})^2\frac{m_1^2+\frac{4}{5}M_W^2\sin^2\beta}{\dfrac{m_1^2}{\cos^2\beta}+\frac{4}{5}M_W^2}\ . \quad (3.7)$$

On the other hand, multiplying the chargino mass matrix from the left with $\mathbf{x}_6=(\cos\varphi,-\sin\varphi)$, with $\tan\varphi=2m_1\sqrt{5}/(4M_W\sin\beta)$, one obtains

$$m_{\tilde{W}_-}^2\le 2M_W^2(1-\sqrt{8/5})^2\frac{m_1^2+\frac{4}{5}M_W^2\cos^2\beta}{\dfrac{m_1^2}{\sin^2\beta}+\frac{4}{5}M_W^2}\ . \quad (3.8)$$

The two bounds (3.7) and (3.8) are obviously identical for $|\tan\beta|=1$, yielding

$$m_{\tilde{W}_-}\le M_W(\sqrt{8/5}-1)\simeq 22 \text{ GeV}\ .$$

If $|\tan\beta|\neq 1$, one of the bounds rises monotonically with m_1^2, while the other one falls with m_1^2. The absolute upper bound is then again realized in the situation where the two bounds are identical, which is true for $m_1^2=0.4M_W^2\sin(2\beta)$; this gives

$$m_{\tilde{W}_-}\le M_W(\sqrt{8/5}-1)\sqrt{|\sin(2\beta)|}$$
$$[\mu_2=0.8M_W^2\sin(2\beta)/m_1]\ . \quad (3.9)$$

The bound (3.9) combined with the PETRA limit[8] $m_{\tilde{W}_-}\gtrsim 22$ GeV already rules out the case when Eq. (2.5) is exactly satisfied unless $|\tan\beta|$ is very close to 1. It also implies that in this region of parameter space the decay $Z\to\tilde{W}_-\tilde{W}_-$ is possible for quite substantial values of $m_{\tilde{Z}_1}$. From Fig. 2(c) we read off that $m_{\tilde{Z}_1}$ has to be larger than about 20 GeV for this decay channel to be kinematically closed.

Unfortunately we have not been able to derive an analytic bound on $m_{\tilde{Z}_2}$ for arbitrary values of $\tan\beta$ for $\tilde{Z}_1$ given by Eq. (2.6). We have, however, numerically checked that $\tan\beta=1$ gives the maximal values of $m_{\tilde{Z}_2}$. In this case $-2m_1$ is an eigenvalue of the neutralino mass matrix (2.1). The remaining two eigenvalues $E_{1,2}$ are then solutions of the equation

$$E^2-E\left[\mu_2(1+\tfrac{5}{3}\tan^2\theta_W)+\frac{8M_W^2}{5\mu_2}\right]+\tfrac{5}{3}\tan^2\theta_W\mu_2^2$$
$$+1.6M_W^2(1+\tfrac{5}{3}\tan^2\theta_W)-M_Z^2=0\ , \quad (3.10)$$

where Eqs. (2.2) and (2.5) have been used. The smaller solution of Eq. (3.10) increases monotonically with μ_2, while according to Eq. (2.5) $2m_1$ decreases with μ_2. The mass of the $\tilde{Z}_2$ is thus maximal if the smaller solution of Eq. (3.10) equals $2m_1=1.6M_W^2/\mu_2$. This is true for $\mu_2=\sqrt{0.6}M_W/\sin\theta_W$, which gives the bound

$$m_{\tilde{Z}_2}\le 8M_W\sin\theta_W/\sqrt{15}\simeq 80.4 \text{ GeV}\quad (2m_1=1.6M_W^2/\mu_2)\ . \quad (3.11)$$

From Fig. 2(c) one sees that this bound depends only weakly on β. Thus $\tilde{Z}_2$ can be substantially heavier than in the cases where the $\tilde{Z}_1$ is a light photinolike or Higgsino-like state. In contrast, the bound on $m_{\tilde{W}_-}$ is much stronger here than in these previously discussed cases.

This can be used to derive nontrivial lower bounds on $m_{\tilde{Z}_2}$ in this region of parameter space, $m_1\gtrsim 0.8M_W^2\sin(2\beta)/\mu_2$, since all sets of parameters that yield $m_{\tilde{W}_-}<22$ GeV are ruled out.[8] The resulting bounds are represented by the dashed curves in Fig. 2(c). For $11\lesssim m_{\tilde{Z}_1}\le 16$ GeV this bound is reached for very large values of $2m_1$, i.e., $\tilde{Z}_1\simeq\lambda_0$ and $\tilde{Z}_2\simeq\lambda_3$, so that one finds (independent of β)

$$m_{\tilde{Z}_2}\simeq m_{\tilde{W}_-}\simeq|\mu_2|\simeq 0.6m_{\tilde{Z}_1}/\tan^2\theta_W\quad (|2m_1|\gg M_Z)\ . \quad (3.12)$$

For $m_{\tilde{Z}_1}<11$ GeV this choice of $2m_1$ would lead to an unacceptably small $m_{\tilde{W}_-}$; the existence of a very light $\tilde{Z}_1$ of the type shown in Eq. (2.6) thus forces the $\tilde{Z}_2$ to be substantially heavier than the value given by Eq. (3.12). For $m_{\tilde{Z}_1}\gtrsim 16$ GeV a new solution occurs with small values of $|2m_1|$ and very large $|\mu_3|$ ($|\mu_3|\gtrsim 3$ TeV). For $m_{\tilde{Z}_1}\ge 22$ GeV this solution corresponds to the limit $|\mu_3|\to\infty$, $m_{\tilde{Z}_1}=m_{\tilde{Z}_2}=m_{\tilde{W}_-}=|2m_1|$, which means that there is no lower bound on $m_{\tilde{Z}_2}$. On the other hand, $|\mu_3|$ should not substantially exceed ~ 1 TeV since otherwise the stability of the gauge hierarchy is no longer guaranteed. We have, therefore, also shown the lower bound on $m_{\tilde{Z}_2}$ with the requirement $|\mu_3|\le 2$ TeV. In this case the bound for $m_{\tilde{Z}_1}\gtrsim 11$ GeV is always reached for $|2m_1|\simeq 1$ TeV; see Eq. (3.12). In contrast, no lower bound on $m_{\tilde{Z}_2}$ can be derived if the $\tilde{Z}_1$ is a light photino or Higgsino.

It is interesting to note that in most cases $m_{\tilde{Z}_2}$ is maximal if it is equal to $m_{\tilde{Z}_3}$; the exception being the case where the $\tilde{Z}_1$ is a light photinolike state with $|\tan\beta|\neq 1$. This means that if $m_{\tilde{Z}_2}$ is close to its upper bound, the model usually predicts $m_{\tilde{Z}_3}$ to be of order M_Z too. However, in general, no strong bounds can be derived for $m_{\tilde{Z}_3}$ or $m_{\tilde{W}_+}$.

IV. GAUGE-BOSON DECAYS INTO CHARGINOS AND NEUTRALINOS

We have seen in the last section that whenever the $\tilde{Z}_1$ is light the decay $Z\to\tilde{Z}_1\tilde{Z}_2$ is kinematically allowed. In most cases the decay $W\to\tilde{W}_-\tilde{Z}_1$ is also possible, the exception being the case where $|\tan\beta|\simeq 1$ and $|2m_1|$ and $|\mu_2|$ are both small. Finally, it may well be possible that even the decays $Z\to\tilde{W}_-\tilde{W}_-$ and $W\to\tilde{W}_-\tilde{Z}_2$ are allowed. In this section we give decay widths for these decays as well as the $Z\to\tilde{Z}_1\tilde{Z}_1$ decay, which is trivially allowed if the $\tilde{Z}_1$ is light.

The relevant couplings for the $Z\to\tilde{W}_-\tilde{W}_-$ and $W\to\tilde{W}_-\tilde{Z}_i$ decays can be found in Refs. 27 and 17, respectively. For the $Z\tilde{Z}_i\tilde{Z}_j$ interactions we find

$$\mathcal{L}_{Z\tilde{Z}_i\tilde{Z}_j}=W_{ij}\bar{\tilde{Z}}_i\gamma_\mu(\gamma_5)^{1+\theta_i+\theta_j}\tilde{Z}_jZ^\mu , \qquad (4.1)$$

where $\theta_i=0$ (1) if the corresponding eigenvalue of the neutralino mass matrix (2.1) is positive (negative). The couplings W_{ij} are given by

$$W_{ij}=\tfrac{1}{4}\sqrt{g^2+g'^2}(i)^{\theta_j}(-i)^{\theta_i}(v_1^{(i)}v_1^{(j)}-v_2^{(i)}v_2^{(j)}) , \qquad (4.2)$$

where the $v_k^{(i)}$ have been defined in the paragraph following Eq. (2.3). These couplings can also be found in the review by Haber and Kane,[1] but in a somewhat different notation.

The decay width for the decay of a vector boson V into two fermions f_1 and f_2 with vector coupling a and axial-vector coupling b is given by

$$\Gamma(V\to f_1f_2)=\frac{\Delta_{12}}{24\pi M_V^3}\lambda^{1/2}(M_V^2,m_{f_1}^2,m_{f_2}^2)\times\left[(a^2+b^2)\left[2M_V^2-m_{f_1}^2-m_{f_2}^2-\frac{(m_{f_1}^2-m_{f_2}^2)^2}{M_V^2}\right]+6(a^2-b^2)m_{f_1}m_{f_2}\right] , \qquad (4.3)$$

where the λ function is as usually defined by

$$\lambda(x,y,z)=x^2+y^2+z^2-2xy-2xz-2yz . \qquad (4.4)$$

The statistics factor Δ_{12} in Eq. (4.3) is $\frac{1}{2}$ if f_1 and f_2 are identical Majorana fermions and 1 otherwise. If f_1 and f_2 are *both* Majorana fermions, the width (4.3) is to be multiplied by another factor of 4.

Note that the expression in the square brackets in Eq. (4.3) vanishes at the edge of phase space, $M_W=m_{f_1}+m_{f_2}$, if $a=0$ and $b\neq 0$, while it approaches the constant $12a^2m_{f_1}m_{f_2}$ if $a\neq 0$ and $b=0$. From Eq. (4.1) we see that the $Z\to\tilde{Z}_i\tilde{Z}_i$ decays always have $a=0$, while the $Z\to\tilde{Z}_1\tilde{Z}_2$ may be mediated by a pure vector interaction if $\theta_1+\theta_2=1$. Thus, the latter decay is favored if the coupling strength (4.2) and the kinematic factors are the same.

It is worth mentioning that most of the $Z\tilde{Z}_i\tilde{Z}_j$ couplings of Eq. (4.2) vanish if $\tan\beta=1$. The reason is that in this case one eigenstate of the neutralino mass matrix (2.1) is the pure Higgsino state with

$$\mathbf{v}^{(\tilde{h})}=(1/\sqrt{2},1/\sqrt{2},0,0) ;$$

obviously there is no $Z\tilde{h}\,\tilde{h}$ coupling. Since the three other eigenvectors have to be orthogonal to the Higgsino the only nonvanishing couplings in this case are the $Z\tilde{h}\,\tilde{Z}_i$ couplings where the $\tilde{Z}_i$ is not the Higgsino state. In our explicit examples we, therefore, avoid the exceptional point $\tan\beta=1$.

At this point it seems appropriate to briefly list the signals that emerge from the decays under consideration. Obviously the $Z\to\tilde{Z}_1\tilde{Z}_1$ decay can only be detected in neutrino counting experiments,[28] since the $\tilde{Z}_1$ escapes direct observation. The $\tilde{Z}_1\tilde{Z}_2$ final state leads to missing transverse momentum ($\not{p}_T$) plus jet(s) or two acoplanar leptons and thus might lead to a striking signature at e^+e^- colliders;[29] if the $\tilde{Z}_2$ decays into $\nu_l\bar{\nu}\tilde{Z}_1$, which might be its dominant decay mode in certain superstring-inspired models,[18] this mode again might be detectable only in neutrino counting experiments. Correspondingly the $Z\to\tilde{Z}_2\tilde{Z}_2$ decay might lead to final states with up to four charged leptons and/or up to four jets. If the $\tilde{Z}_2$-$\tilde{W}_-$ mass splitting is large and squarks and sleptons are heavy compared to M_W, the $\tilde{Z}_2$ dominantly decays[18] into the $\tilde{W}_-$ and two quarks or leptons for $|\tan\beta|\simeq 1$. The $\tilde{W}_-$, in turn, decays into the $\tilde{Z}_1$ plus quarks or a $l\bar{\nu}_l$ pair. The most promising signal from the decay $W\to\tilde{W}_-\tilde{Z}_1$ at a hadron collider comes from jet(s) plus $\not{p}_T$ events; the leptonic signal is swamped by the $W\to l\nu_l$ background.[20] Finally, the $W\to\tilde{W}_-\tilde{Z}_2$ decay leads to final states with three charged leptons and little hadronic activity, or one or two leptons plus jet(s)$+\not{p}_T$.

Unfortunately, the branching ratios for the $\tilde{Z}_2$ and $\tilde{W}_-$ decays strongly depend on the masses of the scalar particles of the model. We therefore do not attempt to quantify our discussion of the signals. Instead we refer the reader to the existing literature[18,20,27] where various cases are discussed in detail.

We are now in a position to analyze the regions of parameter space that lead to a light $\tilde{Z}_1$. We begin with a discussion of the case where the $\tilde{Z}_1$ is photinolike, i.e., the case of small $|\mu_2|$. In Fig. 3 we show the W and Z decay widths as a function of $m_{\tilde{Z}_1}$ for $\tan\beta=2.5$, $2m_1=150$ GeV and negative μ_2; the corresponding values of $m_{\tilde{Z}_2}$ and $m_{\tilde{W}_-}$ are shown on the scales above the figure. Here and in Figs. 4 and 5 the W widths are in units of $\Gamma(W\to e\nu)\simeq 250$ MeV and the Z widths are

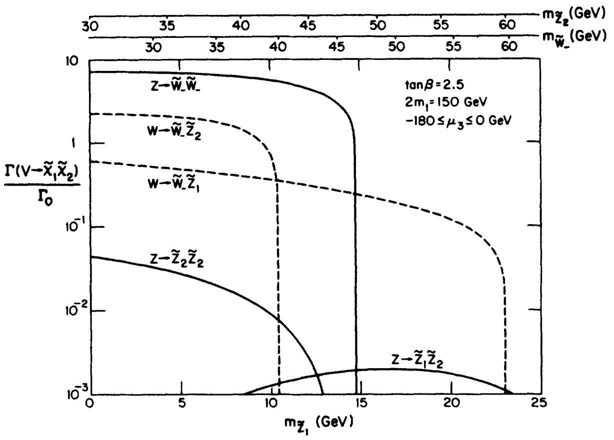

FIG. 3. Decay widths of W bosons (dashed) and Z bosons (solid curves) into neutralinos and charginos as a function of $m_{\tilde{Z}_1}$ for $\tan\beta=2.5$, $2m_1=150$ GeV, and $-180\ \mathrm{GeV}\le\mu_3\le0$, corresponding to the case of a photinolike $\tilde{Z}_1$. The W widths are in units of $\Gamma(W\to e\nu)\simeq250$ MeV, and the Z widths are in units of $\Gamma(Z\to e^+e^-)\simeq92$ MeV. The scales above the frame show the values of $m_{\tilde{W}_-}$ and $m_{\tilde{Z}_2}$ corresponding to $m_{\tilde{Z}_1}$ on the bottom scale.

in units of $\Gamma(Z\to\mu^+\mu^-)\simeq92$ MeV.

Since $\tilde{Z}_1$ is dominantly a photino here it couples only very weakly to the Z boson so that the $Z\to\tilde{Z}_1\tilde{Z}_1$ decay is not observable. If the decay $Z\to\tilde{Z}_2\tilde{Z}_2$ is allowed it leads to clean signatures at SLC and LEP. At LEP I, where $\gtrsim10^5$ $Z\to\mu^+\mu^-$ events are expected annually, there are up to 4000 $\tilde{Z}_2\tilde{Z}_2$ events per year. Because of the small branching fraction ($\lesssim0.006\%$) the decay $Z\to\tilde{Z}_1\tilde{Z}_2$ is at best marginally observable even at these Z^0 factories. It should also be pointed out that the exact rates for the latter two modes strongly decrease if $|2m_1|$ is much larger or if $|\tan\beta|$ approaches one.

In contrast, the $Z\to\tilde{W}_-\tilde{W}_-$ and $W\to\tilde{W}_-\tilde{Z}_i$, $i=1,2$, decay widths are rather insensitive to the exact values of the parameters as long as $|\mu_2|$ is not too large. The $Z\to\tilde{W}_-\tilde{W}_-$ width would, for example, only rise by about 30% if $|2m_1|$ approaches infinity. If open, this decay will thus lead to a substantial increase of the total width of the Z boson.[30] In the region of small $m_{\tilde{Z}_1}$,

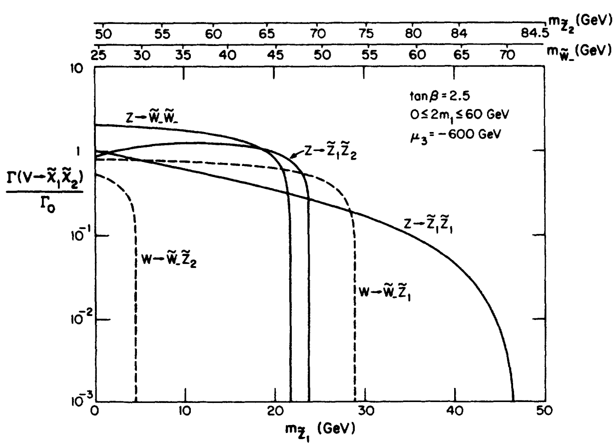

FIG. 4. The same as Fig. 3, but for $0\le2m_1\le60$ GeV and $\mu_3=-600$ GeV, corresponding to the case where $\tilde{Z}_1$ is dominantly a Higgsino. Note that for the given value of μ_3, $m_{\tilde{Z}_2}$ never exceeds 84.5 GeV.

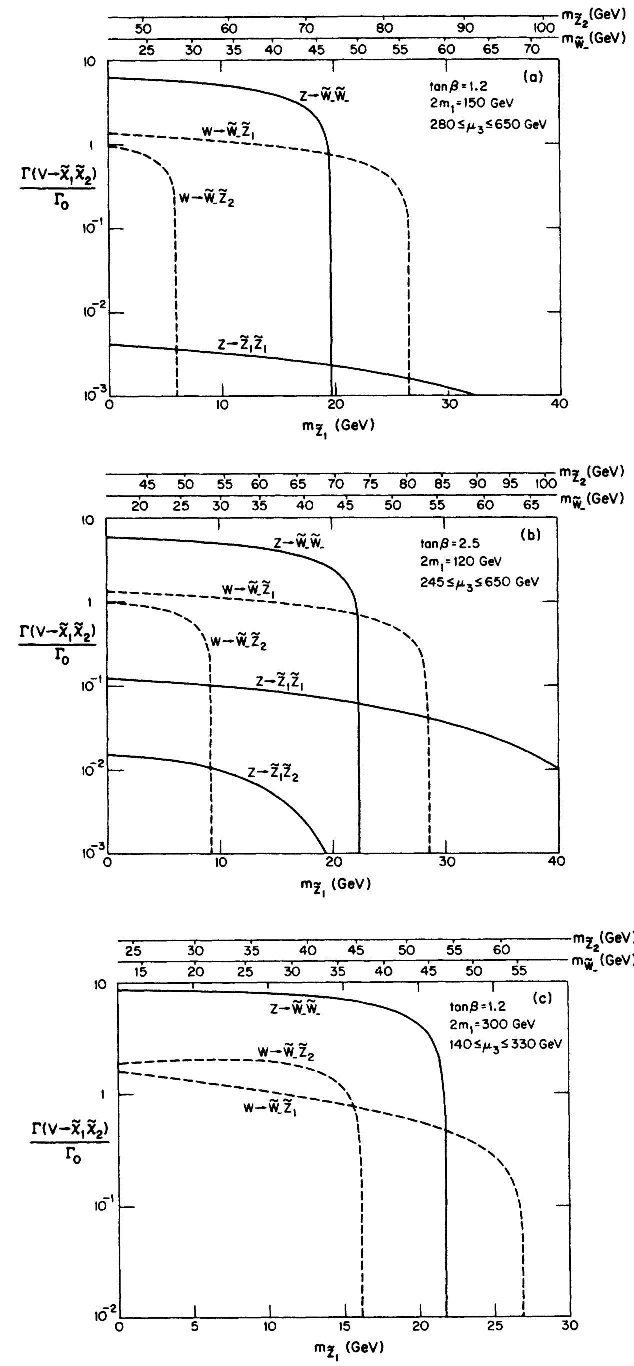

FIG. 5. The same as Fig. 3, but for (a) $\tan\beta=1.2$, $2m_1=150$ GeV, and 280 GeV $\le\mu_3\le650$ GeV, (b) $\tan\beta=2.5$, $2m_1=120$ GeV, and 245 GeV $\le\mu_3\le650$ GeV and (c) $\tan\beta=1.2$, $2m_1=300$ GeV, and 140 GeV $\le\mu_3\le330$ GeV. In all cases $\tilde{Z}_1$ is approximately given by the state (2.6).

which corresponds to $m_{\tilde{W}_-} \lesssim 35$ GeV, the $W \to \tilde{W}_- \tilde{Z}_2$ decay increases the W width by a similar amount. However, this region of parameter space is probably already ruled out by the absence of hadron-free multilepton events in the UA1 data sample.[20] In this case the model predicts an increase of the ratio Γ_Z/Γ_W, unless the $Z \to \tilde{W}_- \tilde{W}_-$ channel is closed.

It has recently been pointed out[31,32] that this might be in conflict with existing data. Combining data from several experiments, the authors of Ref. 32 have concluded that the measured central value of Γ_Z/Γ_W is even smaller than the standard-model prediction with three generations. If this indeed turns out to be the case, the region with $m_{\tilde{W}_-} \lesssim 45$ GeV is disfavored in the case of small μ_2.

The situation changes quite dramatically if the $\tilde{Z}_1$ is dominantly a Higgsino. This is shown in Fig. 4. Note first the relatively large splitting between the masses of the $\tilde{Z}_2$ and the $\tilde{W}_-$; we should mention, however, that this is partly due to our choice of a rather large $|\mu_3|$ ($\mu_3 = -600$ GeV), which is made to ensure that $\tilde{Z}_1$ is dominantly a Higgsino state. Here we find only a small change in Γ_Z/Γ_W compared to the standard-model prediction. Note, however, that substantial contributions to Γ_Z now come from $Z \to \tilde{Z}_1 \tilde{Z}_i$, $i = 1,2$, as well as from $Z \to \tilde{W}_- \tilde{W}_-$. The smallness of this latter decay width can be explained by the observation that the $\tilde{W}_-$ is dominated by its Higgsino component, i.e., the $\tilde{W}_-$ is dominantly a SU(2) doublet rather than a triplet and so has a smaller coupling to Z^0. It is interesting to note that although the $\tilde{W}_-$ and $\tilde{Z}_1$ eigenstates are very different from those of Fig. 3 the $W \to \tilde{W}_- \tilde{Z}_1$ decay width is comparable in the two cases. Finally, the decays $W \to \tilde{W}_- \tilde{Z}_2$ and $Z \to \tilde{Z}_2 \tilde{Z}_2$ are suppressed here since the $\tilde{Z}_2$ is rather heavy ($m_{\tilde{Z}_2} \gtrsim 50$ GeV). At hadron colliders the signals for $\tilde{W}_-$, $\tilde{Z}_1$, and $\tilde{Z}_2$ production from gauge-boson decays are much smaller in this case than in the photino case so that the bounds of Refs. 20 are considerably weakened.

Finally, we turn to a discussion of the case where the $\tilde{Z}_1$ eigenstate is approximately given by Eq. (2.6). Since this region of parameter space has not yet been studied in the existing literature we show plots for three different sets of parameters: $\tan\beta = 1.2$, $2m_1 = 150$ GeV [Fig. 5(a)]; $\tan\beta = 2.5$, $2m_1 = 120$ GeV [Fig. 5(b)]; and $\tan\beta = 1.2$, $2m_1 = 300$ GeV [Fig. 5(c)]. In all the three curves the covered range of μ_3 starts at the value given by Eqs. (2.5) and (2.2), corresponding to $m_{\tilde{Z}_1} = 0$, and ends where all decay modes becomes undetectable.

As indicated by Fig. 2(c), we again observe a rather large difference between $m_{\tilde{Z}_2}$ and $m_{\tilde{W}_-}$, especially in Figs. 5(a) and 5(b) where $2m_1$ is not very large. In these two cases the channel $W \to \tilde{W}_- \tilde{Z}_2$ is already closed at $m_{\tilde{W}_-} = 30$ GeV, which severely limits the possibility of deriving bounds on the parameters from the absence of multilepton $+\not{p}_T$ events at the CERN $Sp\bar{p}S$. In fact, one can see from Fig. 1 that for smaller values of $2m_1$ this decay channel is already closed if $m_{\tilde{W}_-} \gtrsim 22$ GeV, which we know to be true from PETRA experiments.[8] For the values of $2m_1$ shown in Figs. 5(a) and (b), even if this channel is open it contributes only about half as much as in the case of a photinolike $\tilde{Z}_1$.

It is also interesting to note that in this region of parameter space the $Z \to \tilde{W}_- \tilde{W}_-$ decay width is somewhat smaller than in the case of a light photinolike $\tilde{Z}_1$, while the $W \to \tilde{W}_- \tilde{Z}_1$ is about twice as large. Therefore, here the increase in Γ_Z/Γ_W for small values of $m_{\tilde{W}_-}$ is considerably smaller than in the case of a light photino and the same $m_{\tilde{W}_-}$. For the parameters of Fig. 5(a) for $m_{\tilde{W}_-} = 30$ GeV this increase of Γ_Z/Γ_W equals that of about $\frac{2}{3}$ massless neutrinos and almost vanishes at $m_{\tilde{W}_-} = 40$ GeV. (Recall the W and Z decay widths are normalized to 250 and 92 MeV, respectively.)

The comparison of Figs. 5(a) and 5(b) shows that increasing $\tan\beta$ from 1.2 to 2.5 leaves the $Z \to \tilde{W}_- \tilde{W}_-$ and $W \to \tilde{W}_- \tilde{Z}_i$, $i = 1,2$, decay widths almost unchanged, while the widths for $Z \to \tilde{Z}_1 \tilde{Z}_i$ increase by a factor of 30. In the latter case the $Z \to \tilde{Z}_1 \tilde{Z}_2$ decay might be observable at SLC and LEP, while the $\tilde{Z}_1 \tilde{Z}_1$ channel still only counts as at most $\frac{1}{20}$ of a neutrino and will therefore be hard to detect.

On the other hand, the situation for a larger value of $|2m_1|$ depicted in Fig. 5(c) more closely resembles the case where the $\tilde{Z}_1$ is photinolike and $|\tan\beta| \simeq 1$. In the allowed region[8] $m_{\tilde{W}_-} \gtrsim 22$ GeV the difference between $m_{\tilde{Z}_2}$ and $m_{\tilde{W}_-}$ is less than 10 GeV, and the $Z \to \tilde{W}_- \tilde{W}_-$ and $W \to \tilde{W}_- \tilde{Z}_2$ widths are large if these decays are kinematically possible. The reason for these similarities is that the large value of $2m_1$ results in $\tilde{Z}_1$, $\tilde{Z}_2$, and $\tilde{W}_-$ all being dominated by their gaugino components, just as in the case of small $|\mu_2|$ and not too small $2m_1$. As far as gauge-boson decays are concerned, the main difference between the two scenarios is that the case of a light photinolike $\tilde{Z}_1$ leads to a $W \to \tilde{W}_- \tilde{Z}_1$ decay width, which is only about half as large as in the case where the $\tilde{Z}_1$ is approximately given by Eq. (2.6).

V. SCALAR-ELECTRON MASS LIMIT FROM $e^+e^- \to \tilde{Z}_1 \tilde{Z}_1 \gamma$

In the last section we have seen that in most regions of parameter space the $Z \to \tilde{Z}_1 \tilde{Z}_1$ decay width is too small to be observable, even if the $\tilde{Z}_1$ is very light, the exception being the case where the $\tilde{Z}_1$ is dominantly a Higgsino and $|\tan\beta| \neq 1$. This makes it all the more important to look for other reactions in which the LSP is produced together with other visible particles.

At energies well below the Z threshold, i.e., for e^+e^- colliders operating at present, the obvious candidate[33,34] is the reaction $e^+e^- \to \tilde{Z}_1 \tilde{Z}_1 \gamma$. All existing analyses[33,34] have been performed under the simplified assumption that the $\tilde{Z}_1$ is a pure photino $\tilde{\gamma}$. In this case $\tilde{Z}_1$ pairs are produced by t- and u-channel selectron exchange and the photon is radiated off either the beam electrons or the exchanged selectron. With this assumptions the ASP Collaboration has been able to rule out[16] a large region of the $m_{\tilde{e}}-m_{\tilde{Z}_1}$ plane. In Sec. II we have, however, seen that the $\tilde{Z}_1$ can only be a pure photino if $m_{\tilde{Z}_1}=0$ and that even in this extreme case it need not be a photino. In this section we reexamine the existing ASP bound[16] taking into account all the mixings, with particular emphasis on the state (2.6).

The differential cross section for this reaction can be written as

$$d\sigma(e^+e^-\rightarrow\tilde{Z}_1\tilde{Z}_1\gamma)=\frac{e^2|g_{e\tilde{e}\tilde{Z}_1}|^4}{256(2\pi)^4}\left[1-\frac{s'}{s}\right]\left[1-\frac{4m_{\tilde{Z}_1}^2}{s}\right]^{1/2}$$
$$\times\left[\sum_i T_i(P_{e^-},P_{e^+},P_\gamma,K_1,K_2,m_{\tilde{e}},m_{\tilde{Z}_1})\right]d\cos\theta_\gamma d\cos\theta^*_{\tilde{Z}_1}d\phi^*_{\tilde{Z}_1}dx\ . \quad (5.1)$$

Here K_1 and K_2 are the four-momenta of the two $\tilde{Z}_1$, $\theta^*_{\tilde{Z}_1}$ and $\phi^*_{\tilde{Z}_1}$ are the polar and azimuthal angles between them in the $\tilde{Z}_1$-$\tilde{Z}_1$ center-of-mass frame, and θ_γ is the angle between the emitted photon and the beam. Finally, $s'\equiv(K_1+K_2)^2$ and $x\equiv\sqrt{s}/2E_\gamma$. Explicit expressions for the T_i can be found in Ref. 33. To obtain the total cross section the contributions of left- and right-handed electrons have to be added, each of which is given by Eq. (5.1) with the appropriate couplings $g_{e_{L,R}\tilde{e}_{L,R}\tilde{Z}_1}$ and mass $m_{\tilde{e}_{L,R}}$; the interference terms between the contributions of left- and right-handed selectron exchanges are proportional to the electron mass and can thus be ignored. The masses $m_{\tilde{e}_L}$ and $m_{\tilde{e}_R}$ are in general independent parameters. Finally, the $e\tilde{e}\tilde{Z}_1$ couplings are given by[18]

$$g_{e_L\tilde{e}_L\tilde{Z}_1}=(-i)^{\theta_1+1}(1/\sqrt{2})(gv_3^{(1)}+g'v_4^{(1)})\ , \quad (5.2a)$$

$$g_{e_R\tilde{e}_R\tilde{Z}_1}=-(i)^{\theta_1-1}\sqrt{2}g'v_4^{(1)}\ . \quad (5.2b)$$

If the $\tilde{Z}_1$ is a pure photino, Eqs. (5.2) reduce to the known result

$$g_{e_L\tilde{e}_L\tilde{Z}_1}=g_{e_R\tilde{e}_R\tilde{Z}_1}=\sqrt{2}e\ .$$

In order to extract the bound on the cross section we have integrated Eq. (5.1) for the case $\tilde{Z}_1=\tilde{\gamma}$, applying the experimental cuts $E_\gamma\leq 10$ GeV, $p_{T\gamma}>0.8$ GeV, and $20°\leq\theta_\gamma\leq 160°$, and using values of $m_{\tilde{e}}$ and $m_{\tilde{Z}_1}$ that lie on the boundary of the region excluded by the ASP experiment.[16] Using this procedure we found that the experimental limit σ_0 on the $e^+e^-\rightarrow\tilde{Z}_1\tilde{Z}_1\gamma$ cross section depends weakly on $m_{\tilde{Z}_1}$ and can be parametrized by

$$\sigma_0(m_{\tilde{Z}_1})=0.059\ \text{pb}-8.2\times10^{-4}\ \text{pb}\frac{m_{\tilde{Z}_1}}{\text{GeV}}\ . \quad (5.3)$$

We are now in a position to derive bounds on the selectron masses as a function of $m_{\tilde{Z}_1}$ for realistic $e\tilde{e}\tilde{Z}_1$ couplings. To this end we insert the couplings (5.2a) and (5.2b) into Eq. (5.1), and adjust $m_{\tilde{e}}$ such that the cross section equals the value given by Eq. (5.3), where we have, of course, applied the same cuts as before. As mentioned earlier, $m_{\tilde{e}_L}$ and $m_{\tilde{e}_R}$ are, in general, independent parameters. We have therefore investigated the three limiting cases: (I) $m_{\tilde{e}_L}\gg m_{\tilde{e}_R}$, (II) $m_{\tilde{e}_R}\gg m_{\tilde{e}_L}$, and (III) $m_{\tilde{e}_L}=m_{\tilde{e}_R}$. In the first two cases, which give identical bounds on the selectron mass if $\tilde{Z}_1=\tilde{\gamma}$, we have only kept the contributions from the lighter selectron.

We begin with a discussion of the region of small $|\mu_2|$, where the $\tilde{Z}_1$ is similar (though not equal) to the photino. Since σ depends on the fourth power of $g_{e\tilde{e}\tilde{Z}_1}$, even in this region we expect it to substantially deviate from the case where $\tilde{Z}_1=\tilde{\gamma}$, up to 50% for $m_{\tilde{Z}_1}=10$ GeV. However, σ depends equally strongly on the selectron mass. The resulting bounds on $m_{\tilde{e}_L}$ and $m_{\tilde{e}_R}$ therefore differ only slightly from those derived by the ASP Collaboration under the simplified assumption $\tilde{Z}_1=\tilde{\gamma}$.

This is demonstrated in Fig. 6. Here and in Fig. 7 the dashed curves are the original ASP bounds taken from the first paper of Ref. 16. The upper curves are for case (III), $m_{\tilde{e}_L}=m_{\tilde{e}_R}$, while the lower curves have been derived under the assumption that only one selectron-type contributes. Note that for the present case $\tan\beta 2m_1\mu_3<0$ one finds $|g_{e_R\tilde{e}_R\tilde{Z}_1}|>\sqrt{2}e$, $|g_{e_L\tilde{e}_L\tilde{Z}_1}|<\sqrt{2}e$, while for the case $\tan\beta 2m_1\mu_3>0$ (and small $|\mu_2|$) the couplings of the left-handed electrons are enhanced and that of the right-handed electrons suppressed. In both cases these effects tend to cancel if $m_{\tilde{e}_L}\simeq m_{\tilde{e}_R}$, so that the deviations from the original ASP result[16] are even smaller in case (III) than in cases (I) and (II). These results are insensitive to changes in $2m_1$ and $\tan\beta$ as long as μ_2 is sufficiently small.

We note here that we can neglect the Z^0 exchange at SLAC storage ring at PEP beam energies. This can be seen from the fact that the ASP Collaboration can only limit the number of massless neutrino species to less than 7.5 (9.7) at the 90% (95%) confidence level even though these have "full" couplings to Z^0. Thus the con-

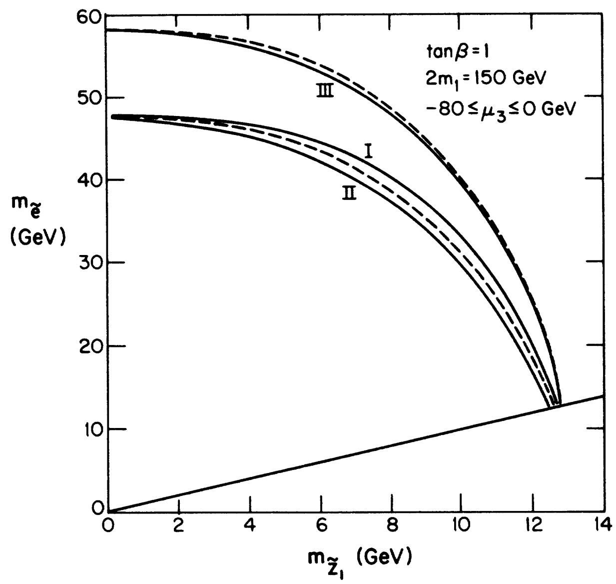

FIG. 6. Excluded region in the $m_{\tilde{Z}_1}$-$m_{\tilde{e}}$ plane for $\tan\beta=1$, $2m_1=150$ GeV, and -80 GeV $\leq\mu_3\leq 0$, corresponding to the case of a photinolike $\tilde{Z}_1$. The dashed curves show the original results[17] of the ASP Collaboration, which have been derived under the assumption that $\tilde{Z}_1$ is exactly a photino; the upper curves are for $m_{\tilde{e}_L}=m_{\tilde{e}_R}$, while the lower curves are for $m_{\tilde{e}_L}\gg m_{\tilde{e}_R}$ or $m_{\tilde{e}_R}\gg m_{\tilde{e}_L}$. For the full curves the neutralino mass matrix (2.1) has been diagonalized exactly, with (I) $m_{\tilde{e}_L}\gg m_{\tilde{e}_R}$, (II) $m_{\tilde{e}_R}\gg m_{\tilde{e}_L}$, and (III) $m_{\tilde{e}_L}=m_{\tilde{e}_R}$. The region below the solid line, $m_{\tilde{e}}<m_{\tilde{Z}_1}$, cannot be excluded by this data set since here the $\tilde{Z}_1$ is not stable; there are, however, much stronger bounds from PETRA experiments on an unstable $\tilde{Z}_1$ which would decay via $\tilde{Z}_1\to\gamma+\not{p}_T$.

tribution from Z^0 exchange to $\tilde{Z}_1$ pair production at PEP is unobservably small for any possible $\tilde{Z}_1$ state.

We now turn to the case where $\tilde{Z}_1$ is dominantly a Higgsino, i.e., to the region of small $|2m_1|$. In this case no bounds can be derived from the ASP data. Even for large values of $|\tan\beta|$ and $m_{\tilde{Z}_1}\simeq 10$ GeV where the gaugino components of $\tilde{Z}_1$ are sizable ($|v_3^{(1)}|,|v_4^{(1)}|\sim 0.15$) the $\tilde{Z}_1\tilde{Z}_1\gamma$ cross section is smaller than σ_0 even for very light selectrons.

We finally turn to the case where the $\tilde{Z}_1$ is approximately given by Eq. (2.6). Some results are shown in Figs. 7(a) and 7(b). Note that $v_3^{(1)}v_4^{(1)}<0$ here, which leads to a least partial cancellation in the $e_L\tilde{e}_L\tilde{Z}_1$ coupling of Eq. (5.2a). If this cancellation is significant, the ASP experiment can obviously not derive any limits on $m_{\tilde{e}_L}$. (Recall that the cross section varies as $g^4_{e_L\tilde{e}_L\tilde{Z}_1}$.)

Figure 7 also shows that the $e_R\tilde{e}_R\tilde{Z}_1$ couplings is suppressed, too, at least if $2m_1$ is not very large, since in this case the $\tilde{Z}_1$ has a sizable Higgsino component which does not contribute to the couplings (5.2b). As expected from Eq. (2.6) the resulting bound on $m_{\tilde{e}_R}$ is rather insensitive to changes of $\tan\beta$. The difference in the results of Figs. 7(a) and 7(b) can be explained by the obser-

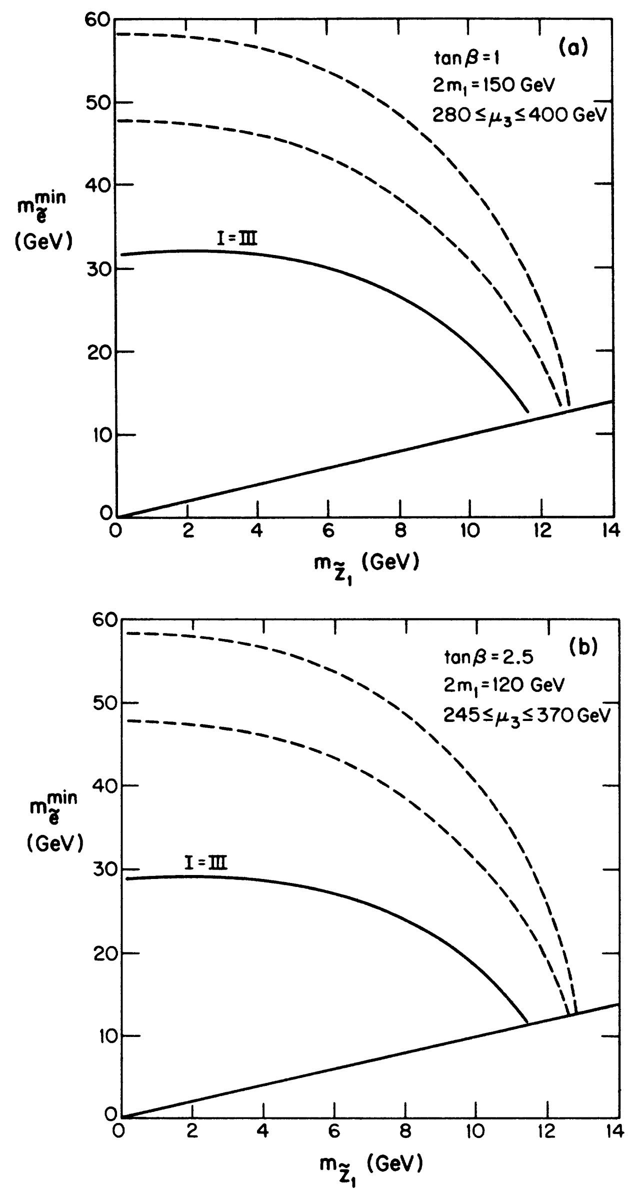

FIG. 7. The same as Fig. 6, but for (a) $\tan\beta=1$, $2m_1=150$ GeV, and 280 GeV $\leq\mu_3\leq 400$ GeV and (b) $\tan\beta=2.5$, $2m_1=120$ GeV, and 245 GeV $\leq\mu_2\leq 370$ GeV. In both cases $\tilde{Z}_1$ is approximately given by the state (2.6). Note that due to the smallness of $g_{e_L\tilde{e}_L\tilde{Z}_1}$, Eq. (5.2a), no limits on $m_{\tilde{e}_L}$ can be derived from the ASP experiment even for $m_{\tilde{Z}_1}=0$.

vation that the smaller value of $2m_1$ of Fig. 7(b) according to Eqs. (2.5) and (2.6) leads to a smaller $|v_4^{(1)}|$. We see from Fig. 7 that if the LSP is approximately given by the state (2.6), the analysis of the ASP experiment is substantially altered, and the limits on selectron masses for fixed $m_{\tilde{Z}_1}$ are very difficult from the $\tilde{Z}_1=\tilde{\gamma}$ case.

VI. THE NEUTRALINO SECTOR OF SUPERSTRING-INSPIRED MODELS AND THE LSP

As is well known, compactification to four dimensions of the anomaly-free ten-dimensional superstring theory[2] is believed to lead to an E(6) supergravity theory[3] with the chiral fields transforming as the **27** representation of the gauge group. E(6) is broken down via the Hosotani mechanism[35] to the low-energy group $G \supset G_S \equiv SU(3)\times SU(2)_L \times U(1)_Y$ with rank $G \geq 5$. These models can be broadly grouped into two classes.[36] In the first group of models[10] G has rank 5 so that it is spontaneously broken to $SU(3)\times U(1)_{em}$ in a single step. In the second class of models G has rank 6. It is first broken to a rank-5 group at an intermediate scale[36,13] $M_I \sim 10^{11}$ GeV and then to G_S at a scale ~ 1 TeV. Here, we ignore the possibility that the last step of the gauge symmetry breaking is due to the vacuum expectation value (VEV) of a "survivor" field,[10] since it has been shown[37] that at least part of the breaking has to be due to VEV's of fields in the generational **27**'s. We further assume for simplicity that the "survivor" fields present[36] in rank-6 models with intermediate scales decouple from the generational **27**'s. At the very least, three Higgs fields are required to break the rank-five subgroup of G to $SU(3)\times U(1)_{em}$. Two of these are the $SU(2)_L$-doublet fields H and $\bar{H}$ that also occur in minimal supergravity models (see Sec. II) while the third is an $SU(2)_L$-singlet field N whose VEV is responsible for the breaking of the rank-5 group to G_S. The main difference between the effective low-energy theories in the case of a rank-5 and a rank-6 group G is the importance of higher-dimensional operators in the effective superpotential in the latter case. This can change the spectrum of the theory in a crucial way as we will see later.

Following our earlier discussion, we will assume that the LSP is in the neutralino sector. In the case of superstring-inspired models, this sector consists of the nine fields H_i, $\bar{H}_i$, and N_i (here $i=1-3$ denotes the generation) and the three neutral gaugino fields for the SU(2), U(1), and extra U(1) gauginos. We work in a basis where only the VEV's of H_3, $\bar{H}_3$, and N_3 do not vanish. Since we are dealing with the low-energy theory, the additional gaugino and Higgsino fields for rank-6 models do not enter our consideration as these have masses $\sim M_I$ (Ref. 36). In general, the twelve neutral fermion fields all mix in a complicated fashion to yield the mass eigenstates, one of which (by assumption) is the LSP.

For rank-5 models or the class of rank-6 models where the breaking at the intermediate scale is due to a VEV of a right-handed sneutrino field, it has been shown that there are always two states (other than $\tilde{\gamma}$) whose masses are smaller than $\sim M_W$ and which may be considerably smaller.[15] One of these states may even be the LSP. At this point we note that if the intermediate scale breaking is due to the VEV of an N-type field, because of higher-dimensional operators in the superpotential the 12×12 neutral mass matrix under discussion gets additional diagonal entries $\sim M_I^2/M_{\text{Planck}} \sim 1$ TeV, so that the considerations of Ref. 15 no longer apply. In this case the LSP is dominantly an $SU(2)_L\times U(1)$ gaugino[10] and the phenomenology is much the same as discussed for supergravity models. In the remainder of this section we focus on the possibility that there is an additional light state, which we denote by $\tilde{n}$ (which is dominantly a combination of N_1 and N_2 if all superpotential couplings are comparable), over and above a relatively light gaugino-like state, which we denote by $\tilde{x}$. $\tilde{x}$ may be approximately the photino, but for larger SUSY-breaking gaugino masses is more the $U(1)_Y$ gauge fermion. $\tilde{x}$ and $\tilde{n}$ decay into one another depending on their masses.

The phenomenology of these decays is very model dependent since the couplings of $\tilde{n}$ to matter are completely unknown. Our considerations naturally divide into two parts: (i) $\tilde{n}$ has no large couplings to matter and (ii) the couplings of $\tilde{n}$ to matter are sizable. These can arise from either the gaugino components of $\tilde{n}$ or, more importantly, from the components of those combinations of H ($\bar{H}$) orthogonal to the combination that develops a VEV. (Recall only the latter has couplings proportional to fermion masses.)

In the first case, $\tilde{x}$ can decay into $\tilde{n}f\bar{f}$ (or the other way if $m_{\tilde{n}} > m_{\tilde{x}}$) via virtual Z or sfermion exchanges. The $\tilde{x}\tilde{n}Z$ coupling comes from the $SU(2)_L$-doublet Higgsino components of both $\tilde{x}$ and $\tilde{n}$ and is thus suppressed by $\sim (M_W/M_{Z'})^2 \sim 10^{-2}$. This suppression follows from the structure of the mass matrix. Similarly, the $f\tilde{f}\tilde{n}$ vertex, which comes from the gaugino content of $\tilde{n}$, is suppressed by the same factor. Therefore, if the scalar fermions are substantially heavier than M_Z, the decays $\tilde{x}\to\tilde{n}f\bar{f}$ (or $\tilde{n}\to\tilde{x}f\bar{f}$) dominantly occur via Z^0 exchange and hence the branchings into the various modes are the same as for the decays of Z^0. If $m_{\tilde{x}} > m_{\tilde{n}}$, this would make it difficult to use $\not{p}_T$ as a footprint of SUSY at hadron colliders since squarks and gluinos would decay into $\tilde{x}$ which would then cascade into $\tilde{n}$ + quarks or leptons. The signatures would instead be similar to the signatures from the decays of squarks (gluinos) into charginos and neutralinos.[17] In case (ii) with sizable superpotential couplings of $\tilde{n}$ to matter, the scalar-fermion-exchange graphs may dominate the Z^0 exchange graph and so the branching fractions are completely governed by these couplings. If, for example, $\tilde{n}$ couples only to leptons, this would produce very interesting events from the decay of scalar quarks which would have to first decay to $\tilde{x}$ via $\tilde{q}\to q\tilde{x}$ followed by $\tilde{x}\to\tilde{n}l\bar{l}$.

One special case that has been extensively considered in the literature[10,13–15,18] is when the 12×12 neutral fermion matrix breaks up into two 6×6 matrices with the first block containing the gauginos and the fermionic partners of the three scalar fields that develop VEV's and the other block the rest. This may[38] or may not[14] be a consequence of a discrete symmetry. The decays of $\tilde{x}$ ($\tilde{n}$) depend on whether or not there is such a symmetry. In the generic case where there is no discrete symmetry, there are matter couplings to $\tilde{n}$ so that $\tilde{x}$ decays again occur via virtual sfermion exchange. If ν^c is essentially massless as in rank-5 models,[10] the decay $\tilde{x}\to\tilde{n}\nu\nu^c$ is also possible. If neither of these decays are possible, $\tilde{x}$ decays via exotic charge $-\frac{1}{3}$ quark loops or via charged $H_{1,2}$ loops into $\gamma+\tilde{n}$. It is worth pointing out that in the case the neutralino sector is split into two sectors there is no coupling of the gauge bosons to fermions in the two different sectors so that the decays $\tilde{x} \to \tilde{n} + \gamma$ via gauge loops[19] are absent. For the same reason, there is no $W \to \tilde{W} \tilde{n}$ decay even if it is kinematically allowed. Finally, if there is a discrete symmetry as considered in Ref. 38, $\tilde{x}$ and $\tilde{n}$ carry different quantum numbers under this symmetry so that $\tilde{x} \to \tilde{n}$ + ordinary particles is forbidden. In this case, the only allowed mode is $\tilde{x} \to \nu \nu^c \tilde{n}$. If ν^c is heavy, $\tilde{x}$ is stable. Also, $\tilde{n}$ cannot couple to usual quarks and leptons so that the phenomenology of the LSP in the ordinary sector is effectively the same as in supergravity models.

Finally, we remark that if there is a second light state $\tilde{n}_2$ in the neutralino sector (recall from our earlier discussion that there may be two light states), it would decay into the lighter of these and a fermion pair via a virtual Z^0 exchange (or via scalar exchange if there were couplings to matter), since the $Z^0 \tilde{n}_1 \tilde{n}_2$ is essentially a gauge coupling up to doublet mixing angles.

VII. SUMMARY AND CONCLUDING REMARKS

With the assumption that the LSP is a state in the neutralino sector we have studied various candidates for the LSP. For the minimal supergravity model we find that in addition to the well-studied cases of the photino and the Higgsino, there is yet one more possibility for a light neutralino state [see Eqs. (2.5) and (2.6)] that may even be massless depending on the parameters of the model. In fact, as can be seen in Figs. 1, there is a relatively large region of parameter space in which the LSP is just this state. A comparison of the resulting phenomenology depending on the nature of the LSP forms the main subject of this paper.

It has been shown[26] using very general arguments that if SUSY-breaking gaugino masses are small so that the LSP is essentially a light photino, the lightest state ($\tilde{W}_-$) in the chargino sector and the next lightest state ($\tilde{Z}_2$) in the neutralino sector have masses smaller than M_W and M_Z, respectively. For the minimal supergravity model we show (Sec. III) that this bound for $m_{\tilde{Z}_2}$ is even more restrictive, $m_{\tilde{Z}_2} < M_Z/\sqrt{2}$. We have further shown that $\tilde{W}_-$ and $\tilde{Z}_2$ are lighter than the W and Z bosons even if the gaugino masses are large provided only that the lightest neutralino ($\tilde{Z}_1$) has a mass $\ll M_W$. The explicit bounds depend on the nature of $\tilde{Z}_1$ and are summarized in Eqs. (3.2), (3.6), and (3.9) and in Figs. 2(a)–2(c) for the three different possibilities for $\tilde{Z}_1$.

The decays of gauge bosons $W^\pm$ and Z^0 into charginos and neutralinos are studied in Sec. IV. The decay rates are shown in Figs. 3–5 for the cases in which the LSP is a photino, a Higgsino, and the state (2.6). We see that except in Fig. 4 the $Z^0 \to \tilde{W}_- \tilde{W}_-$ width is large. This is because $\tilde{W}_-$ is dominantly a Higgsino in this case and so does not have large isotriplet couplings to Z^0. We also see that the rate for the decay $W \to \tilde{W}_- \tilde{Z}_2$ that gives the best signatures for gauginos at the CERN collider is generically large only in the case the LSP is a photino. This would significantly alter the analysis of at least the multilepton final state performed in Ref. 20, although qualitative results can probably be extracted (except for the Higgsino case) by scaling the widths in Figs. 5 with those in Fig. 3 since the kinematics of the reaction does not depend on the nature of the LSP. For the case where the Higgsino is the LSP, the best signals come from $Z \to \tilde{Z}_1 \tilde{Z}_2$ and $Z \to \tilde{W}_- \tilde{W}_-$. The decays of $\tilde{Z}_2$ would lead to very distinctive events particularly at e^+e^- colliders such as SLC or LEP. Yet another important distinction between the photino and the other cases is the rather large splitting between the $\tilde{W}_-$ and $\tilde{Z}_2$ masses in the latter cases [except in Fig. 5(c) where $2m_1$ is large]. This also explains why the $W \to \tilde{W}_- \tilde{Z}_2$ decay width is small except in Fig. 3. Finally, we note that the difference in the decay widths of the vector bosons into charginos and neutralinos would effect the ratio Γ_W/Γ_Z in a model-dependent way.

The mass limits from the ASP experiment[16] are reexamined taking into account all the mixings appropriate to minimal supergravity models. For the case when $\tilde{Z}_1 \approx \tilde{\gamma}$ the change in the excluded region from the pure photino case (shown as the dashed lines) is shown in Fig. 6 for the extreme cases $m_{\tilde{e}_L} \gg m_{\tilde{e}_R}$, $m_{\tilde{e}_L} \gg m_{\tilde{e}_R}$, and $m_{\tilde{e}_L} = m_{\tilde{e}_R}$. We see that the limit on the selectron mass (for a fixed $m_{\tilde{Z}_1}$) is not very different from the pure photino limit quoted in Ref. 16. The corresponding situation when $\tilde{Z}_1$ is the state (2.6) is completely different, as can be seen in Figs. 7(a) and 7(b). We see that even for a massless $\tilde{Z}_1$ the limit on the right-handed selectron mass is at most ~ 30 GeV as compared with almost 50–60 GeV in the photino case. In fact, for the case where $m_{\tilde{e}_R}$ is too heavy to contribute to $\tilde{Z}_1$ pair production, no limit can be deduced on $m_{\tilde{e}_L}$. This is due to the cancellation between the SU(2) and U(1) gaugino contributions to the $e\tilde{e}_L \tilde{Z}_1$ couplings in Eq. (5.2).

Finally, we have qualitatively analyzed the phenomenology of the LSP in a nonminimal supergravity model, using for illustrative purposes a superstring-inspired E(6) model. We find that the nature of the LSP, and hence the resulting phenomenology, is very model dependent. In addition to the three candidates for the LSP discussed earlier, there is the additional possibility that the LSP is an exotic fermion, $\tilde{n}$, i.e., a combination of $SU(2)_L$-doublet and -singlet fields in the **27** of E(6) that have the internal quantum numbers of Higgs bosons but whose scalar partners do not develop a VEV. We have shown that if $\tilde{n}$ is indeed the LSP, there are at least four classes of models (depending on the symmetry of the superpotential), each with rather different phenomenology as discussed in Sec. VI.

To summarize, we have shown that the LSP even in minimal supergravity models need not be the photino or the Higgsino, but may be a more complicated combination of gaugino and Higgsino states. If this is the case, the resulting phenomenology can be substantially altered from the LSP$\approx$photino case usually studied in the literature so that some of the signals for supersymmetry may merit further examination.

ACKNOWLEDGMENTS

We are grateful to R. Godbole for her comments. This research was supported in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation, and in part by the U.S. Department of Energy under Contract No. DE-AC02-76ER00881.

[1]H. E. Haber and G. L. Kane, Phys. Rep. **117**, 75 (1985); S. Dawson, E. Eichten, and C. Quigg, Phys. Rev. D **31**, 1581 (1985); P. Nath, R. Arnowitt, and A. Chamseddine, *Applied* $N=1$ *Supergravity,* ICTP Series in Theoretical Physics (World Scientific, Singapore, 1984), Vol. I; H. P. Nilles, Phys. Rep. **110**, 1 (1984).

[2]M. Green and J. Schwarz, Phys. Lett. **149B**, 117 (1984).

[3]P. Candelas, G. Horowitz, A. Strominger, and E. Witten, Nucl. Phys. **B258**, 46 (1985); E. Witten, *ibid.* **B258**, 75 (1985).

[4]S. Wolfram, Phys. Lett. **82B**, 65 (1979); C. B. Dover, T. K. Gaisser, and G. Steigman, Phys. Rev. Lett. **42**, 1117 (1979); D. Dicus and V. Teplitz, *ibid.* **44**, 218 (1980).

[5]P. F. Smith and J. R. J. Bennett, Nucl. Phys. **B149**, 525 (1979); P. F. Smith *et al.*, *ibid.* **B206**, 333 (1982), and references cited therein; E. B. Norman, S. B. Gazes, and D. A. Bennett, Phys. Rev. Lett. **58**, 1403 (1987).

[6]R. Arnowitt, A. Chamseddine, and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); **50**, 232 (1983); L. Ibáñez, Phys. Lett. **118B**, 73 (1982); R. Barbieri, S. Ferrara, and C. Savoy, *ibid.* **119B**, 343 (1982); J. Ellis, L. Ibáñez, and G. G. Ross, Nucl. Phys. **B221**, 29 (1983); L. E. Ibáñez and C. Lopez, *ibid.* **B233**, 511 (1984); C. Kounnas, A. B. Lahanas, D. V. Nanopoulos, and M. Quiros, *ibid.* **B236**, 438 (1984); L. E. Ibáñez, D. Lopez, and C. Muñoz, *ibid.* **B256**, 218 (1985).

[7]J. S. Hagelin, G. L. Kane, and S. Raby, Nucl. Phys. **B247**, 638 (1984); L. Ibáñez, Phys. Lett. **137B**, 160 (1984); R. M. Barnett, H. E. Haber, and K. S. Lackner, *ibid.* **126B**, 64 (1983); Phys. Rev. D **29**, 1990 (1984).

[8]See, for example, M. Davier, in *Proceedings of the XXIII International Conference on High Energy Physics,* Berkeley, California, 1986, edited by S. C. Loken (World Scientific, Singapore, 1987).

[9]I. J. Kroll, in *Electroweak Interactions and Unified Theories,* proceedings of the XXII Rencontre de Moriond, Les Arcs, France, 1987, edited by J. Tran Thanh Van (World Scientific, Singapore, 1987); F. Pauss, Proceedings of the Workshop on Experiments, Detectors and Experimental Areas for the Supercollider, Berkeley, 1987 (unpublished).

[10]J. Ellis, K. Enqvist, D. Nanopoulos, and F. Zwirner, Nucl. Phys. **B276**, 14 (1986).

[11]P. Nath, R. Arnowitt, and A. Chamseddine, Phys. Lett. **129B**, 445 (1983; D. Dicus, S. Nandi, and X. Tata, *ibid.* **129B**, 451 (1983); V. Barger *et al.*, *ibid.* **131B**, 372 (1983); J. Ellis, *et al.*, *ibid.* **132B**, 436 (1983); L. Alvarez-Guame, J. Polchinski, and M. Wise, Nucl. Phys. **B221**, 495 (1983). These five papers are examples of early work on models with a light photino. For upgrade see Ref. 1 and references therein.

[12]H. E. Haber, G. L. Kane, and M. Quiros, Phys. Lett. **160B**, 297 (1985); Nucl. Phys. **B273**, 333 (1986).

[13]L. Ibáñez and J. Mas, Nucl. Phys. **B286**, 107 (1987).

[14]J. Ellis, D. Nanopoulos, S. Petcov, and F. Zwirner, Nucl. Phys. **B283**, 93 (1987).

[15]M. Drees and X. Tata, Phys. Lett. B **196**, 65 (1987).

[16]C. Hearty *et al.*, Phys. Rev. Lett. **58**, 1711 (1987); see also G. Bartha *et al.*, *ibid.* **56**, 685 (1986); E. Fernandez *et al.*, *ibid.* **54**, 1118 (1985).

[17]H. Baer, V. Barger, D. Karatas, and X. Tata, Phys. Rev. D **36**, 96 (1987). We note that the doublet fields are denoted in a different notation. h and h' of this reference are $\bar{h}$ and h, respectively, in the present paper.

[18]H. Baer, D. Dicus, M. Drees, and X. Tata, Phys. Rev. D **36**, 1363 (1987).

[19]H. Komatsu and J. Kubo, Phys. Lett. **157B**, 90 (1985); Nucl. Phys. **B263**, 265 (1986).

[20]H. Baer, K. Hagiwara, and X. Tata, Phys. Rev. Lett. **57**, 294 (1986); Phys. Rev. D **35**, 1598 (1987); A. Chamseddine, P. Nath, and R. Arnowitt, Phys. Lett. B **174**, 399 (1986).

[21]M. Glück, R. Godbole, and E. Reya, Phys. Lett. B **186**, 421 (1987).

[22]C. Aljabar *et al.*, Phys. Lett. B **185**, 241 (1987).

[23]V. Barger, N. Deshpande, and K. Whisnant, Phys. Rev. Lett. **56**, 30 (1986); S. M. Barr, *ibid.* **55**, 2778 (1985); E. Cohen, K. Enqvist, J. Ellis, and D. V. Nanopoulos, Phys. Lett. **165B**, 76 (1985); M. Drees, N. Falck, and M. Glück, *ibid.* **167B**, 187 (1986); M. Durkin and P. Langacker, *ibid.* **166B**, 436 (1986); U. Amaldi *et al.*, Phys. Rev. D **36**, 1385 (1987). A recent analysis by V. Barger and K. Whisnant *ibid.* **36**, 3429 (1987)] suggests that for the model of Ellis, Enqvist, Nanopoulos, and Zwirner (Ref. 10), the extra Z boson is heavier than 1 TeV.

[24]G. Steigman, K. Olive, D. Schramm, and M. Turner, Phys. Lett. B **176**, 33 (1986); J. Ellis, K. Enqvist, D. Nanopoulos, and S. Sarkar, Phys. Lett. **167B**, 452 (1986).

[25]M. Drees, (Nucl. Phys. B, in press).

[26]S. Weinberg, Phys. Rev. Lett. **50**, 387 (1983); Arnowitt, Chamseddine, and Nath (Ref. 6).

[27]X. Tata and D. Dicus, Phys. Rev. D **35**, 2110 (1987).

[28]E. Ma and J. Okada, Phys. Rev. Lett. **41**, 287 (1978); K. J. F. Gaemers, R. Gastmans and F. M. Renard, Phys. Rev. D **19**, 1605 (1979).

[29]D. Dicus, S. Nandi, W. Repko, and X. Tata, Phys. Rev. D **29**, 1317 (1984); **30**, 1112 (1985); H. Baer *et al.*, in *Physics at LEP,* LEP Jamboree, Geneva, Switzerland, 1985, edited by J. Ellis and R. Peccei (CERN Report No. 86-02, Geneva, 1986); H. Komatsu, Phys. Lett. B **177**, 201 (1986).

[30]V. Barger, R. Robinett, W.-Y. Keung, and R. Phillips, Phys. Rev. D **28**, 2912 (1983); D. Dicus, S. Nandi, W. Repko, and X. Tata, Phys. Rev. Lett. **50**, 1030 (1983).

[31]A. Datta and R. Godbole, Phys. Rev. D **37**, 225 (1988); A. Martin, R. G. Roberts, and W. J. Stirling, Phys. Lett. B **189**, 220 (1987).

[32]F. Halzen, C. S. Kim, and S. S. D. Willenbrock, Phys. Rev. D **37**, 229 (1988).

[33]K. Grassie and P. N. Pandita, Phys. Rev. D **30**, 22 (1984).

[34]This process was first considered by P. Fayet, Phys. Lett. **117B**, 460 (1982); and has also been subsequently considered by J. Ellis and J. Hagelin, *ibid.* **122B**, 303 (1983); T. Kobayashi and M. Kuroda, *ibid.* **139B**, 208 (1984); J. Ware

and M. Machacek, *ibid*. **142B**, 300 (1984).
[35]Y. Hosotani, Phys. Lett. **129B**, 193 (1983).
[36]M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi, and N. Seiberg, Nucl. Phys. **B259**, 519 (1985).
[37]P. Kalyniak and M. Sundaresan, Phys. Lett. **167B**, 320 (1986).
[38]M. Drees and X. Tata, Phys. Rev. Lett. **59**, 528 (1987).