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How model dependent are sparticle mass bounds from CERN LEP?

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In order to translate nonobservation of events of a type not expected in the standard model into bounds on particle masses, one necessarily has to make some specific assumptions. We systematically study the model dependence of the bounds that can be derived from the negative outcome of sparticle and Higgs-boson searches at CERN LEP as well as from the measurement of the total and nonhadronic decay widths of the Z boson. We also show that the constraints from direct searches taken together with those from the width measurements imply stronger bounds than those obtained using either strategy by itself. Regarding model dependence, we find that the bounds on charged-sparticle and sneutrino masses that have been obtained with the usual assumptions for their decay patterns are, in fact, generally valid within the framework of the minimal supersymmetric model (MSSM). In models with a conserved R parity, these bounds can be significantly reduced only by the *ad hoc* introduction of two charged SU(2)-singlet and/or an additional pair of SU(2)-doublet superfields. The bounds on neutralino masses, on the other hand, are very sensitive to even minor changes of the minimal model. At present, even in the MSSM the mass of the lightest neutralino is only mildly constrained and, in fact, cannot be bounded in models with just a slightly altered neutralino sector unless a supersymmetry signal is found. We also comment on the implications of a possible negative outcome of sparticle searches in future experiments.

I. INTRODUCTION

Supersymmetry[1] (SUSY) is a very attractive extension of the well-established standard model (SM) of strong and electroweak interactions. In spite of the fact that there is no evidence for any deviation from the SM in the data either from the increased luminosity runs at the CERN S$p\bar{p}$S and the Fermilab Tevatron or from the decay properties of the Z^0 boson derived from the sample of $\sim 10^5$ Z^0's obtained last year at the SLAC Linear Collider (SLC) or CERN LEP experiments, the unnaturalness of the SM symmetry-breaking sector strongly suggests the existence of new physics at a scale $\Lambda \lesssim 1$ TeV. In particular, the motivation for searching for sparticles with $m \lesssim \Lambda$ remains strong especially because model calculations suggest that the lighter sparticle masses may be as small as few tens of GeV provided the heaviest of these have masses $\leq \Lambda$. The lack of experimental evidence for any physics beyond the SM is, of course, interpreted as a lower limit on sparticle masses.

These mass bounds are necessarily model dependent. First, the decay of any sparticle, and hence its experimental signature, depends on the properties of the daughters. Since in many models sparticles only decay into other sparticles the signature of any sparticle depends on the properties of other sparticles lighter than the parent.

In most of the analyses that lead to bounds on sparticle masses from the collider data, the production rate is only dependent on the gauge couplings of the sparticles and so is fixed by their SU(3)×SU(2)×U(1) quantum numbers (this is what makes SUSY somewhat predictive in the first place). In practice, of course, this is complicated by the fact that the sparticles are model-dependent mixtures of states with definite gauge quantum numbers.

These ambiguities exist even within the so-called minimal supersymmetric model[1] (MSSM), which is essentially a direct supersymmetrization of the SM; well-known examples are the dependence[2] of bounds on the masses of strongly interacting sparticles on details of the electroweak gaugino-Higgsino sector, and the dependence[3] of the $Z^0 \to$ charginos decay width on the relative magnitude of the gaugino and Higgsino components of the chargino.

In this paper, we systematically study the model dependence of bounds on sparticle masses that have recently been obtained from the study of Z^0 decays at SLC and LEP.[4] Where appropriate, we also comment on future search strategies. We restrict ourselves to models with a conserved R parity, where sparticles can only be produced in pairs and the lightest supersymmetric particle (LSP) is stable and escapes detection. We focus on this case since it is already well known[5] that R-parity-violating models can lead to very different yet distinctive experimental signatures since the LSP is unstable. Since a (stable) LSP has to be[6] electrically and color neutral unless it is very heavy (in which case its mass exceeds $\sim \Lambda$) it can only be (i) the scalar partner of the neutrino, the sneutrino ($\widetilde{\nu}$) (or in extended models a neutral R-odd spin-zero particle), (ii) the neutralino defined as an R-odd spin-$\frac{1}{2}$ neutral sparticle or (iii) the gravitino. We ignore the last of these possibilities since, in this case, the next-lightest sparticle is very long lived (unless the gravitino is ultralight[7]) and so behaves as a canonical LSP in collider experiments provided it is neutral.

The main purpose of this paper is to study the extent to which the bounds on sparticle masses that are obtained from the data on the decays of Z^0, and derived within the framework of the MSSM, are dependent on the model assumptions. To quantify this, we have constructed "minimally nonminimal" models to evade each of these bounds. It is not our aim to advocate any of these models, some of which are rather contrived. It should be stressed that the main argument that favors the MSSM over other supersymmetric models is its simplicity. However, there are other only slightly more complicated models which can be argued to be no less (and sometimes, even more) appealing than the MSSM. For instance, the model with an additional Higgs singlet does not necessitate the introduction of a dimensionful parameter in the superpotential. We show that some bounds like the one on the lightest neutralino, discussed in Sec. II B, can be evaded by minor modifications of the MSSM. We believe that it is important to stress that such bounds should not be considered to be "generic" bounds on supersymmetry. In contrast, we find that many bounds, e.g., the one on the lightest chargino, can only be evaded by adding several extra superfields and/or unnatural tuning of several parameters; these bounds obviously have a much more general validity. However, even these more complicated models should not always be dismissed lightly. For instance, it turns out that the minimal extension of the MSSM (discussed in Sec. III B) necessary to evade the chargino bound also allows for the purely radiative generation of neutrino masses, and hence perhaps for a natural solution to the solar-neutrino problem. But quite apart from these considerations, we believe it is important to recognize in just what classes of models each of the recently published limits on supersymmetric particle masses is valid.

Apart from the various bounds that have been obtained by directly searching for SUSY signals, we will also make use[8] of the fact that the total and nonhadronic widths of the Z have been measured[9] and found to be in good agreement with the SM. The 90% upper limits on any non-SM contributions to these quantities are

$$\delta\Gamma(Z,\text{total}) < 96\ \text{MeV}, \tag{1.1}$$

$$\delta\Gamma(Z,\text{nonhadronic}) < 38.3\ \text{MeV}. \tag{1.2}$$

The bound (1.2) has been derived from the hadronic peak cross section under the assumption that the hadronic decay width of the Z^0 boson is as given by the SM.

The remainder of this paper is organized as follows. In Sec. II, after a brief recapitulation of the assumptions of the MSSM, we discuss the bounds on sparticle masses resulting from negative results for direct searches for sparticles and the neutral Higgs boson, as well as from the precise measurement of the Z^0 width and the hadronic cross section at the peak. The interplay between the bounds that can be obtained using different methods is especially emphasized. In Sec. III, we discuss various extensions of the neutralino and chargino sectors of the MSSM that could lead to substantial modifications of the bounds obtained in Sec. II. In Sec. IV, we discuss alterations of the MSSM that could potentially modify the bounds on slepton masses. Section V contains a summary of our results and an outlook for the future.

II. THE MINIMAL MODEL

When experimenters quote bounds on sparticle masses, they are usually meant to be valid within some region of parameter space of the MSSM. A precise definition of this model seems in order. Like all "realistic" SUSY models, the MSSM is a field theory with supersymmetry softly broken at the scale Λ (discussed in Sec. I) since otherwise the smallness of M_W is difficult to understand without resorting to fine-tuning.[1] In addition, the minimal model has the following defining properties.

(i) The gauge group relevant for TeV scale physics is $SU(3) \times SU(2) \times U(1)_Y$; i.e., there are no heavy gauge bosons beyond the W and the Z^0.

(ii) There are no new matter superfields besides those containing the quarks and leptons present in the SM.

(iii) The Higgs sector contains just the two doublet superfields needed[10] to give supersymmetric masses to all the quarks and charged leptons and to cancel anomalies.

(iv) The SUSY-breaking Majorana masses of the different gauginos are equal at the unification scale. Since the renormalization-group (RG) scaling of these masses is identical to that of the squared gauge couplings, we have

$$\mu_3 = \alpha_s \frac{\sin^2\theta_W}{\alpha_{\text{em}}}\mu_2 = \frac{3}{5}\frac{\alpha_s}{\alpha_{\text{em}}}\cos^2\theta_W \mu_1, \tag{2.1}$$

where μ_3, μ_2, and μ_1 are the masses of SU(3), SU(2), and U(1) gaugino, respectively.

Conditions (i) and (ii) are just the statement that the matter fermion and gauge boson content are as in the three-generation SM. The significance of (iii) is often not spelled out. First, we note that in supersymmetric theories one cannot introduce new Higgs bosons without introducing their Higgsino partners, which via their mixing with the gauginos and Higgsinos of the MSSM can change the couplings of all the charginos and neutralinos. This, in turn, can change the decay patterns of all unstable sparticles. The additional Higgs bosons, via their gauge interactions, can also mediate decays of charginos and neutralinos, thereby possibly altering the signatures expected from the MSSM. From conditions (i) and (iii) it follows, assuming that R parity is conserved, that the MSSM contains two charginos $\widetilde{W}_-$, $\widetilde{W}_+$ (with $m_{\widetilde{W}_-} \leq m_{\widetilde{W}_+}$) and four neutralinos $\widetilde{Z}_i$, $\widetilde{Z}_1$ being the lightest and $\widetilde{Z}_4$ the heaviest.

Condition (iii) is obviously significant even in the nonsupersymmetric sector. For instance, from the absence of the decays[11,12] $Z \to Z^* H$ and $Z \to HP$, the LEP experiments have been able to significantly constrain the Higgs sector of the MSSM which is determined[13] by two additional parameters, which may be taken to be v'/v the ratio of the vacuum expectation values of the fields, h' and h that give masses to the $T_3 = -\frac{1}{2}$ and $T_3 = +\frac{1}{2}$ fermions, respectively, and one of the scalar masses. Here H and P are the lighter neutral scalar and pseudoscalar, respectively. Since $m_t \gg m_b$, all known models predict $v' < v$. Assuming that this is indeed the case, the absolute bounds from the LEP experiments can be written as

$$M_H > 21 \text{ GeV}, \tag{2.2}$$

$$v'/v < 0.77, \tag{2.3}$$

where it is understood that the bound on the Higgs boson may be significantly stronger than (2.2) depending on v'/v. As we will see, (2.3) has important consequences for slepton and neutralino searches.

Finally, we note that (iv) enables us to translate experimental bounds on the gluino mass $m_{\widetilde{g}} = |\mu_3|$ to those on μ_1 and μ_2 via Eq. (2.1). As stated above, (2.1) follows from the equality of masses (which, admittedly, may be altered by the introduction of complicated gauge kinetic terms) at the unification scale and is independent of the details of the physics between the weak and unification scales. This distinguishes (2.1) from any relation between squark and slepton masses (beyond those imposed by SU(3)×SU(2)×U(1) gauge invariance), which depend differently on the degrees of freedom contributing to their evolution than do the gauge couplings. Lower bounds on the gluino mass in the vicinity of 80 GeV have been obtained[14] from the non-observation of missing transverse momentum events at hadron colliders, assuming the gluino always decays via $\widetilde{g} \to q\,\bar{q}\,\widetilde{Z}_1$. However, this bound can be considerably relaxed[2] if other decays are possible. In obtaining our bounds on μ_2 and μ_1, we will conservatively assume that

$$m_{\widetilde{g}} = |\mu_3| > 50 \text{ GeV} \tag{2.4}$$

We now turn to the discussion of the restriction on the parameters of the MSSM obtained from the lack of any evidence for SUSY in Z^0 decays at LEP, beginning with sleptons.

A. Sleptons

Within the MSSM, there are three different kinds of sleptons: sneutrinos $\widetilde{\nu}$, and the superpartners $\widetilde{l}_L$ and $\widetilde{l}_R$ of left-handed and right-handed charged leptons.

1. Sneutrinos

A light $\widetilde{\nu}$, being neutral, is a candidate for the LSP. If this is the case, or if $\widetilde{Z}_1$ is the only sparticle lighter than $\widetilde{\nu}$, the sneutrino is either stable or decays invisibly via $\widetilde{\nu} \to \nu\,\widetilde{Z}_1$ through either the (usually small) Z-ino component of $\widetilde{Z}$, or via loop diagrams.[15] Thus, sneutrinos lighter than $M_Z/2$ will contribute to the invisible decay width of the Z^0 boson without altering the hadronic decay width from its SM value so that the bound (1.2) is applicable. This then implies[8] that

$$\begin{aligned} &m_{\widetilde{\nu}} \geq 28.9 \text{ GeV(only one light } \widetilde{\nu}), \\ &m_{\widetilde{\nu}} \geq 38.4 \text{ GeV(three degenerate light } \widetilde{\nu}\text{'s)} \end{aligned} \tag{2.5}$$

at 90% C.L.

These bounds are not valid if the $\widetilde{\nu}$ has a sizable hadronic branching ratio. We will see in Sec. II B that within the MSSM the decay $\widetilde{\nu} \to \widetilde{W}_- + l$ is not possible for sneutrinos that violate the bounds (2.5), due to strong lower bounds on $m_{\widetilde{W}_-}$. That leaves the possibility of $\widetilde{\nu} \to \nu + \widetilde{Z}_2$ decays with subsequent hadronic decay of $\widetilde{Z}_2$.

In the limiting case where $\widetilde{Z}_2$ only decays hadronically, sneutrino pair events would contain several jets together with a large amount of $\not{p}_T$ from the escaping neutrinos. For sneutrino masses violating (2.5), the branching fraction for Z^0 decays to $\widetilde{\nu}\widetilde{\nu}$ always exceeds ~1.5 % so that ~ 10^3 such characteristic n jet + $\not{p}_T$ would already be present in the data sample of the four LEP experiments. If $m_{\widetilde{Z}_2} \simeq m_{\widetilde{\nu}}$, so that the neutrinos are very soft, then $\widetilde{\nu}$ - pair events would have the signature of $\widetilde{Z}_2\widetilde{Z}_2$ events but occur at a considerably larger rate. They would thus have been detected by the ALEPH search[16] for neutralinos discussed in the next subsection. We thus conclude the bounds (2.5) are almost certainly applicable; even if $m_{\widetilde{Z}_2} \simeq m_{\widetilde{Z}_1}$, so that the ALEPH bound on $m_{\widetilde{Z}_2}$ does not apply, (2.5) probably does, since, once again sneutrino decays are almost invisible. We note that it is possible that $\widetilde{\nu} \to \nu\widetilde{Z}_2 \to \nu\widetilde{Z}_1 H$, in which case each sneutrino pair event will contain four b quarks.

2. Left-handed charged sleptons

Experimenters at LEP and SLC have searched[4] for charged sleptons, assuming 100% branching ratio for the direct decay into lepton + $\widetilde{Z}_1$. However, within the MSSM, these bounds are already obsolete, as far as left-handed sleptons are concerned. The reason is that the masses of these sleptons are related to the masses of the corresponding sneutrinos:

$$m^2_{\tilde{l}_L} = m^2_{\tilde{\nu}_e} + M^2_W \, \frac{1-(v'/v)^2}{1+(v'/v)^2} \,. \tag{2.6}$$

This is a direct consequence of SU(2) gauge invariance. The resulting mass limit can only be evaded if $\tilde{l}_L$ mixes with some other field to form the mass eigenstate. Note that the bound implied by (2.6) does not apply even if this other field has the same SU(2)×U(1) quantum numbers as $\tilde{l}_L$, so that the coupling of Z^0 to the mass eigenstate pair does not deviate from the $Z^0\,\tilde{l}_L\,\tilde{l}_L$ coupling of the MSSM. If the field that mixes with $\tilde{l}_L$ has different quantum numbers the coupling of the Z^0 to this mass eigenstate is altered so that neither this bound (2.6) nor the direct search bounds apply. This will be discussed in detail in Sec. IV.

Since we have just convinced ourselves that within the MSSM $m_{\tilde{\nu}} > 28.9$ GeV, the bound (2.3) on v'/v implies

$$m_{\tilde{l}_L} > 49.8 \text{ GeV} \tag{2.7}$$

for any left-handed slepton. This bound improves to 55.8 GeV if all sneutrinos are degenerate.

3. Right-handed charged sleptons

Their masses are not related to sneutrino masses. Furthermore, they couple only fairly weakly to Z bosons; even a massless $\tilde{l}_R$ contributes only 17.6 MeV to the Z^0 width. Even the bound (1.2) does therefore only lead to a nontrivial constraint if all three generations are degenerate, in which case it only implies

$$m_{\tilde{l}_R} > 20 \text{ GeV} \qquad (\text{three degenerate } \tilde{l}_R) \tag{2.8}$$

at 90 % C.L. We therefore have to rely on the direct searches.[4] The best bounds to date come from OPAL and exclude the regions $m_{\tilde{e}_R} < 42$ GeV, $m_{\tilde{\mu}_R} < 41$ GeV, and $m_{\tilde{\tau}_R} < 40$ GeV if the slepton always decays into the corresponding lepton and a light, stable neutralino, though the other experiments have comparable bounds. These bounds can be evaded if $m_{\widetilde{Z}_1}$ is close to $m_{\tilde{l}_R}$, since in this case the produced leptons are soft and hence difficult to distinguish from two-photon and beam-gas backgrounds. However, making use of the fact that due to the Z^0 peak the signal would still show a strong dependence on the beam energy while the background is rather flat, the DELPHI Collaboration has been able[17] to shrink the allowed value of $m_{\tilde{l}_R} - m_{\widetilde{Z}_1}$ to a little over one GeV for $\tilde{l} = \tilde{e}, \tilde{\mu}$ and $m_{\tilde{l}_R} \lesssim 38$ GeV. One exception to this analysis may occur if $\tilde{l}_R$ dominantly decays to $\widetilde{Z}_2$. In the MSSM, this can easily occur if $\widetilde{Z}_1 \approx \tilde{h}$ and $\widetilde{Z}_2 \approx \tilde{\gamma}$. Even in this case, an $\tilde{l}_R\,\bar{\tilde{l}}_R$ event ought to be distinctive since it would necessarily contain the two primary (very likely, hard) leptons along with the decay products of $\widetilde{Z}_2$ (even if these are hadronic). Only in the unlikely possibility that $m_{\tilde{\tau}_R} \ll m_{\tilde{e},\tilde{\mu}}$ is it at all conceivable that $Z \to \bar{\tilde{\tau}}_R \tilde{\tau}_R$ events (whose signature will be two τ's accompanied by up to two more tau pairs or other hadronic debris) may be missed. Even so, the event would look rather unusual in that it would have a very high "jet" multiplicity (if the τ's are confused with jets), high sphericity and abnormally narrow jets.

B. Charginos and neutralinos

Signatures for chargino and neutralino production in Z^0 decays have been extensively studied[3,18,19] within the framework of the MSSM. The masses and couplings of all the charginos and neutralinos relevant for our analysis are fixed in terms of just three parameters. In the notation of Ref. 19, we may take these to be the gluino mass ($m_{\tilde{g}}$), the ratio v'/v of the vevs introduced earlier and the supersymmetric Higgsino mixing mass $2m_1$. Apart from naturalness and fine-tuning arguments which suggest that all the mass parameters are $\lesssim 1$ TeV and that $v'/v \sim 1$ and the experimental limits on $v\prime/v$ (Ref. 12) and $m_{\tilde{g}}$ (Ref. 14) discussed above, there are no a priori constraints on these parameters.

Most direct searches[4] have focused on the lighter chargino, $\widetilde{W}_-$, since it is expected to be copiously produced in Z^0 decays provided it is kinematically accessible. These searches exclude a region of the $m_{\widetilde{W}_-} - m_{\widetilde{Z}_1}$ plane, under the assumption that the chargino always decays via $\widetilde{W}_- \to \widetilde{Z}_1 f\bar{f}'$ and that this decay is dominated by virtual W exchange (so that, except for phase-space effects which might be important if $m_{\widetilde{W}_-} - m_{\widetilde{Z}_1} \lesssim$ few GeV, the branching fraction into any particular $f\bar{f}'$ mode is just given by the corresponding one for W$\to f\bar{f}'$). This is probably a conservative assumption in that the left-handed squarks which can mediate hadronic decay of $\widetilde{W}$ are known to be at least as heavy as their slepton counterparts and, unless[2] the chargino is only just above the reach of LEP I, heavier[14] than ~ 75 GeV. Thus, sfermion-mediated decays of the chargino, if anything, may enhance its somewhat cleaner leptonic decays from the values predicted by just the W-exchange diagrams.

The Z$\to \widetilde{W}_-\overline{\widetilde{W}}_-$ decay width depends, of course, on the composition of the $\widetilde{W}_-$ mass eigenstate in terms of W-ino and Higgsino current states. Within the MSSM, one finds that a pure W-ino (Higgsino) maximizes (minimizes) this decay width. The ALEPH Collaboration have announced[20] lower limits on the chargino mass for these extreme cases; obviously their limit depends on $m_{\widetilde{Z}_1}$ since for $m_{\widetilde{W}_-} = m_{\widetilde{Z}_1}$ there can be no visible decay products, and hence no limit. Their bounds can approximately be parametrized as follows.

Pure W-ino:

$m_{\widetilde{W}_-} - m_{\widetilde{Z}_1} < \Delta_1\ (m_{\widetilde{W}_-})$, with

$$\Delta_1(m_{\widetilde{W}_-}) = \begin{cases} 0.19 m_{\widetilde{W}_-} - 1\ \text{GeV}, & m_{\widetilde{W}_-} \le 40\ \text{GeV}, \\ m_{\widetilde{W}_-} - 31.4\ \text{GeV}, & 40\ \text{GeV} \le m_{\widetilde{W}_-} \le M_Z/2. \end{cases} \tag{2.9a}$$

Pure Higgsino:

$m_{\widetilde{W}_-} - m_{\widetilde{Z}_1} < \Delta_2\ (m_{\widetilde{W}_-})$, with

$$\Delta_2(m_{\widetilde{W}_-}) = \begin{cases} 0.31 m_{\widetilde{W}_-} - 1.3\ \text{GeV}, & m_{\widetilde{W}_-} \le 40\ \text{GeV}, \\ m_{\widetilde{W}_-} - 28.9\ \text{GeV}, & 40\ \text{GeV} \le m_{\widetilde{W}_-} \le M_Z/2. \end{cases} \tag{2.9b}$$

Since the two extremes can only be realized if either $m_{\widetilde{g}}$ or $|\ 2m_1\ |$ is infinite, a realistic chargino will have both Higgsino and gaugino components. We assume that the bound on $m_{\widetilde{W}_-} - m_{\widetilde{Z}_1}$ for this realistic case, derived by linearly interpolating between (2.9a) and (2.9b) via

$$m_{\widetilde{W}_-} - m_{\widetilde{Z}_1} \le \Delta_2(m_{\widetilde{W}_-}) - \frac{\Gamma_R(m_{\widetilde{W}_-}) - \Gamma_2(m_{\widetilde{W}_-})}{\Gamma_1(m_{\widetilde{W}_-}) - \Gamma_2(m_{\widetilde{W}_-})}\left[\Delta_2(m_{\widetilde{W}_-}) - \Delta_1(m_{\widetilde{W}})\right] \tag{2.10}$$

is valid, where Γ_1, Γ_2, and Γ_R are, respectively, the $Z \to \widetilde{W}_-\overline{\widetilde{W}}_-$ widths for the case of a pure W-ino, a pure Higgsino and the "realistic" chargino.

In order to obtain a bound on $m_{\widetilde{W}_-}$, we have scanned the space of parameters $(v\prime/v, m_{\widetilde{g}}, 2m_1)$ and for each point we have computed the chargino and neutralino masses, the total $Z^0 \to$ gauginos width, and the $Z^0 \to \widetilde{W}_-\overline{\widetilde{W}}_-$ width. We find that the constraint (1.1) on the total Z^0 width alone suffices[8] to derive the bound

$$m_{\widetilde{W}_-} \ge 40\ \text{GeV}. \tag{2.11a}$$

If, in addition, we impose the bound (2.10) from the direct chargino search[20] [since this bound is obtained assuming the chargino can only decay to the LSP, we use (2.10) only when $m_{\widetilde{Z}_2} > m_{\widetilde{W}_-}$], this bounds improves to

$$m_{\widetilde{W}_-} \ge 43\ \text{GeV}, \tag{2.11b}$$

independent of the mass of any other sparticle. This is a good example which illustrates how bounds from inclusive decays of the Z^0 can collaborate with bounds from direct searches resulting in a stronger bound than from either method alone.

Notice also that the constraint (2.11b) within the MSSM framework implies that $m_{\widetilde{W}_-} - m_{\widetilde{Z}_1} > 13$ GeV (5 GeV) provided $m_{\widetilde{g}} < 1$ TeV (3 TeV). This, of course, means that one would not have to worry about the possibility that a produced chargino might escape detection (particularly at an e^+e^- collider) because its visible decay products are too soft. Unfortunately, in non-minimal models (such as the free μ_1 model discussed in the next section), this is no longer true. We now turn to constraints on the MSSM parameters from neutralino decay of the Z^0.

The ALEPH Collaboration,[16] based on their search for $Z \to \widetilde{Z}_1\widetilde{Z}_2$ and $\widetilde{Z}_2\widetilde{Z}_2$ decays has announced limits on neutralino decays of Z^0. In addition to the tree-level decay $\widetilde{Z}_2 \to \widetilde{Z}_1\ H$ or $\widetilde{Z}_1\ f\bar{f}$, they have also included the one loop radiative decay $\widetilde{Z}_2 \to \widetilde{Z}_1\gamma$ in their analysis since this can be substantial[21] for certain ranges of SUSY parameters. Their results can be approximately parametrized as follows.

$\widetilde{Z}_1\widetilde{Z}_2$ decays: (i) if $m_{\widetilde{Z}_1} > 0.9 m_{\widetilde{Z}_2}$ there is no bound; (ii) if $1.2 m_{\widetilde{Z}_2} - 52$ GeV $\le m_{\widetilde{Z}_1} \le 0.5\ m_{\widetilde{Z}_2}$, $B(Z \to \widetilde{Z}_1\widetilde{Z}_2) \le 3 \times 10^{-4}$; (iii) if $m_{\widetilde{Z}_2} = M_Z$ and $m_{\widetilde{Z}_1} = 0$, $B(Z \to \widetilde{Z}_1\widetilde{Z}_2) < 10^{-3}$.

For any set of values for $(m_{\widetilde{Z}_1}, m_{\widetilde{Z}_2})$, we obtain the bound by linear interpolation between the boundary of region (ii) and the line, $m_{\widetilde{Z}_1} = 0.9\ m_{\widetilde{Z}_2}$ or the point $(0, M_Z)$ as the case may be. If d_1 is the distance of $(m_{\widetilde{Z}_1}, m_{\widetilde{Z}_2})$ to this line or point, and d_2 the corresponding distance to the boundary of region (ii), we have by linear interpolation that the branching fraction $B(Z \to \widetilde{Z}_1\widetilde{Z}_2)$ satisfies

$$B(Z \to \widetilde{Z}_1\widetilde{Z}_2) \le \frac{3 \times 10^{-4} d_1 + 10^{-3} d_2}{d_1 + d_2}. \tag{2.12}$$

$\widetilde{Z}_2\ \widetilde{Z}_2$ decays: the region where B$(Z \to \widetilde{Z}_2\widetilde{Z}_2) < 1.1 \times 10^{-3}$, which we conservatively take to approximate the current bound from ALEPH, is given by

$$m_{\widetilde{Z}_1} \ge f(m_{\widetilde{Z}_2})\ , \tag{2.13a}$$

where

$$f(x) = \begin{cases} x, & x \le 8\ \text{GeV}, \\ 0.5x + 4\ \text{GeV}, & 8\ \text{GeV} \le x \le 15\ \text{GeV}, \\ x - 3.5\ \text{GeV}, & 15\ \text{GeV} \le x \le 41\ \text{GeV}, \\ 37.5\ \text{GeV}, & 41\ \text{GeV} \le x \le M_Z/2. \end{cases} \tag{2.13b}$$

In using the constraints (2.12) and (2.13) to obtain bounds on the SUSY parameter space, we have considered the decay of an unpolarized Z^0; i.e., we have ignored spin correlations both due to neutralinos and due to the Z^0. Including these introduces a dependence of the neutralino angular distribution on the sign of the neutralino eigenvalues which can play a role since the experiment does not have 100% angular coverage. Since the ALEPH constraints[16] on the branching fractions have been obtained always using the more conservative case, we believe that our bounds are valid despite this approximation.

We have checked that the constraints (2.12) and (2.13) do not lead to any bound on $m_{\widetilde{Z}_1}$ and $m_{\widetilde{Z}_2}$ [except $m_{\widetilde{Z}_2} \geq 8$ GeV, which follows from the bound (2.4) on the gluino mass]. Here, we have conservatively ignored the production of neutralinos other than $\widetilde{Z}_1$ and $\widetilde{Z}_2$, since the experimental analysis[16] assumes that the unstable neutralino directly decays to $\widetilde{Z}_1$ and fermions, whereas within the MSSM framework, it is quite possible that $\widetilde{Z}_3$ decays via $\widetilde{Z}_2$ which subsequently decays to the lightest neutralino. If, instead of the analysis discussed above, we take the bounds (2.12) and (2.13) to be those on all possible neutralino decays of Z^0 except into the LSP pair, i.e., we assume that the heavy neutralinos all decay directly to the LSP so that the analysis of Ref. 16 is valid, we do find a lower bound of 4 GeV on $m_{\widetilde{Z}_1}$. We have also checked that in the region of parameter space where the contribution of $\widetilde{Z}_3$ is crucial to obtain this bound, $\widetilde{Z}_3$ does, in fact, dominantly decay to $\widetilde{Z}_1$ so that the ALEPH analysis is valid. The existence of such a bound has already been noted by Roszkowski,[22] who has shown that the region of parameter space excluded by ALEPH when combined with the UA2 bound[14] $m_{\tilde{g}} > 79$ GeV implies a lower bound $m_{\widetilde{Z}_1} < 10$–13 GeV. On the other hand, as already discussed in Ref. 8, the constraint (1.1) on the total width leads to an interesting bound on $m_{\widetilde{Z}_3}$. The constraints (2.3) on v'/v and (2.10) on the charginos now strengthen this to

$$m_{\widetilde{Z}_3} \geq 61 \text{ GeV} . \tag{2.14}$$

The bound on the lightest neutralino mass as well as that on $m_{\widetilde{Z}_2}$ can be considerably improved in future experiments. As the Z sample increases, the bound on the invisible decay width of the Z as well as that on the visible $Z \to$ neutralino decays are expected to improve. Furthermore, searches for Higgs bosons can strengthen the bound on v'/v. Some preliminary improvements have been recently reported.[23] The combined upper bound on the invisible decay width of the Z boson is now in the vicinity of 25 MeV. Also, the ALEPH Collaboration has sharpened the bound (2.3) to $v'/v < 0.625$,[24] and has

TABLE I. Present and expected lower bounds on neutralino masses from a study of Z decays at LEP assuming the minimum gluino mass, the largest value of v'/v (as obtained from Higgs-boson searches), and the largest value of the NSM contribution to the invisible Z width are as shown in the first three columns for an "improvement factor" x over the bounds on the visible neutralino decays of the Z as given by (2.12) and (2.13) of the text. The entries labeled $x = 1$ thus correspond to the present published bounds. Also shown in parentheses in the last two columns are the values $(2m_1, m_{\tilde{g}})$ of the SUSY parameters for which the minimum values of the neutralino masses shown in the last two columns are realized. This occurs when v'/v takes on the largest possible value consistent with the Higgs boson constraints.

$m_{\tilde{g},\text{min}}$ (GeV)	v'/v (max)	Γ (MeV)	x[a]	$m_{\widetilde{Z}_1,\text{min}}$ (GeV)	$(m_{\tilde{g}}, 2m_1)$	$m_{\widetilde{Z}_2,\text{min}}$ (GeV)	$(m_{\tilde{g}}, 2m_1)$
50	0.77	38.3	1	4	(50, +4)	8	(50, + 4)
50	0.77	25	1	4	(50, +4)	8	(50, + 4)
50	0.625	25	1	8.1	(50, 125)	16	(95, 17)
50	0.625	25	5	8.1	(50, 125)	17.3	(105, 19)
50	0.625	25	25	8.1	(50, 125)	17.3	(105, 19)
50	0.625	25	125	16.5	(105, 19)	17.3	(105, 19)
100	0.625	25	1	13.5	(158, 15)	17.3	(105, 19)
100	0.625	25	5	16	(103, 264)	17.3	(105, 19)
100	0.625	25	25	16	(103, 264)	17.3	(105, 19)
100	0.625	25	125	16.5	(105, 19)	17.3	(105, 19)
100	0.5	25	5	15.5	(100, 210)	35.8	(220, 33)
100	0.5	25	25	15.5	(100, 210)	37.3	(230, 32)
100	0.5	25	125	20.0	(135, 483)	38.8	(240, 32)
150	0.5	25	5	22	(150, 500)		
150	0.5	25	25	22	(150, 500)	as for	$m_{\tilde{g}} > 100$ GeV
150	0.5	25	125	22	(150, 500)		

[a]Improvement factor over the published ALEPH bound of Ref. 16.

also reported improved bounds from a negative outcome of their neutralino search. Within the framework of models with a common gaugino mass at some high-energy scale, these limits will also depend on the gluino mass limit that would be obtained from hadron colliders.

These are summarized in Table I where we have shown the current limits obtained from the analysis reported above using the constraints (1.1), (1.2), (2.3), (2.9), (2.12), and (2.13) as well as the "improved limits" using the recent results reported in Ref. 23. Also shown by the $x \neq 1$ entries is how much the bound on neutralino masses can be expected to improve with a large enough Z sample so that the limits on the branching fractions are improved by a factor x as compared to the present published bounds, (2.12) and (2.13). The value for $x = 5$ roughly corresponds to the preliminary bounds recently reported by ALEPH.[23] In our analysis, if no limit on $m_{\widetilde{Z}_2}$ is possible from these constraints because it is almost degenerate with $\widetilde{Z}_1$, we have included all its contributions to the Z width in $\Gamma_{\rm inv}$. We have also shown in the table the values of the minimal model parameters for which the bounds on the two lightest neutralino masses are realized. We note the following.

(i) Using the published values of the constraints and treating the $\widetilde{Z}_3$ as discussed above, we find the lightest neutralino bound is just 4 GeV. If an improved limit of about 100 GeV is assumed on the gluino mass, this moves up to a higher value, consistent with Ref. 22. Using the recent preliminary data presented at the Singapore Meeting[23] improves this number to 8 GeV (13.5 GeV) for $m_{\tilde{g}} < 50$ GeV (100 GeV).

(ii) We see that a further increase in the sensitivity to the visible decays of Z^0 leads to only a slight improvement in the neutralino mass bounds unless either the gluino mass bound, the v'/v bound from Higgs-boson searches or the invisible width bound is strengthened.

(iii) It is interesting to note that even with a 125-fold increase in sensitivity (this may well be impossible due to backgrounds), the two lightest neutralinos may be as light as 17 GeV even if $m_{\tilde{g}}$ is as heavy as 100 GeV. Since $\widetilde{Z}_2$ is almost degenerate with $\widetilde{Z}_1$, its decay products are invisible so that this limit can be improved either by increasing the sensitivity with which the invisible width can be measured or (as shown by the $v'/v = 0.5$ entries) if the Higgs sector is further constrained.

(iv) Even for gluino masses close to the reach of the Tevatron and v'/v as small as 0.5 (corresponding to a Higgs boson heavier than 54 GeV), the lightest neutralino of the MSSM need be no heavier than about 22 GeV. Note that this is so for rather natural values of the mass parameters and is independent of the size of the data sample. Thus it is unlikely that these bounds will significantly improve beyond the values in the Table I until the beam energy of LEP is increased. We note that improving the sensitivity of LEP to the invisible decays of the Z does not alleviate the situation. To see this, we note that in the limit of very large $2m_1$ (with $m_{\tilde{g}}$ fixed), the two lightest neutralinos are essentially U(1) and SU(2) gauginos (with masses approximately $m_{\tilde{g}}/7$ and $m_{\tilde{g}}/3.5$, respectively) and so decouple from the Z whereas the heavier neutralinos have masses equal to $2m_1$ and so are kinematically inaccessible in Z decays. We have also checked that for $m_{\tilde{g}} < 150$ GeV (the range that can be accessed at the Tevatron) $2m_1 = 500$ GeV is large enough to suppress the branching fraction for the neutralino decays of the Z to below 10^{-6}. In the same limit, the chargino is a pure SU(2) charged gaugino with a mass, $m_{\tilde{g}}/3.5$, and so has a very large coupling to Z and thus can be readily searched for as soon as the energy of LEP is increased. Until this time (or until a better limit on the gluino mass is available from hadron colliders), we believe that it will be impossible to improve the limit on the neutralino masses beyond the values shown in Table I by a study of Z decays alone.

In the future, it may also be possible to constrain the neutralino masses by a study of W decays. Analyses[25] exist only for W bosons collected prior to 1985 and, since then, the data sample has increased by an order of magnitude in size. As shown in Fig. 1, a measurement of the total width of the W may provide the key, though this may have to wait till LEP II since the width measurement via the determination of $R = N(W \to e\nu)/N(Z \to ee)$ at a hadron collider is unlikely to achieve the necessary precision. In Fig. 1, we show the minimal (solid) and maximal (dashed) width for $W \to \widetilde{W}_- \widetilde{Z}_i (i = 1,2)$ decays after all existing constraints [(1.1), (1.2), (2.3), (2.4), (2.10), (2.12) and (2.13)] have been included. Note that the bound (2.4) on the gluino mass together with the bound (2.11) on $m_{\widetilde{W}_-}$ implies that a very light $\widetilde{Z}_1$ ($m_{\widetilde{Z}_1} \leq 8$ GeV) has to be Higgsino-like. From Fig. 1 we see that such a light Higgsino should have been produced

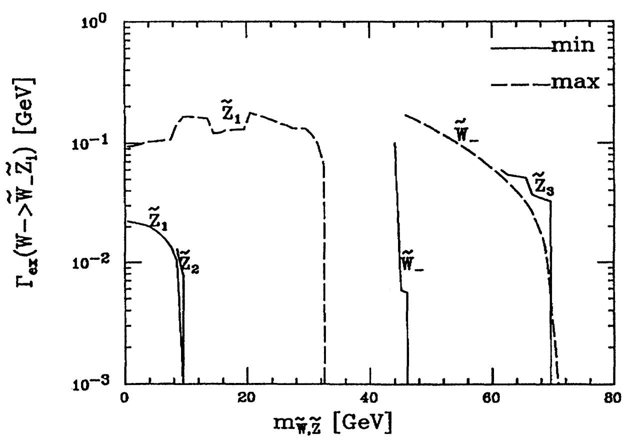

FIG. 1. Minimal (solid) and maximal (dashed) values of $\sum_i \Gamma(W \to \widetilde{W}_- \widetilde{Z}_i)$ in the MSSM as a function of various neutralino and chargino masses. All the LEP constraints discussed in the text have been included; they are responsible for the structure of the maximal width as a function of $m_{\widetilde{Z}_1}$. Notice that for a massless $\widetilde{Z}_1$ the width is determined to be 45 MeV within a factor of 2.

in roughly 0.7% of all W decays or more. The absence of the corresponding signal would imply a lower bound on $m_{\widetilde{Z}_1}$, and hence on the mass of the LSP, within the MSSM; this mass plays an important role, e.g., in cosmology, since the LSP is[26] an excellent dark-matter candidate.

Note furthermore that the supersymmetric decay width of the W boson can still be as large as 150 MeV, even if $m_{\widetilde{Z}_1} \simeq 30$ GeV. We see also that although a chargino as heavy as 70 GeV could be produced in W decays with a branching fraction in excess of 1%, whether it can be detected would depend on its decays. The order-of-magnitude increase in the size of the W sample should lead to an improvement of the monojet bound claimed earlier[25] to > 50 GeV. Also, if we are lucky, and the chargino dominantly decays via $\widetilde{W}_- \to \widetilde{l}\nu$ or $l\widetilde{\nu}$ (which, if $m_{\widetilde{W}_-} \simeq 70$ GeV would be consistent with all LEP constraints), its signal would show up as a large excess in "$W \to \tau \to l$ decays" at hadron colliders and probe regions of SUSY parameter space not accessible at LEP. Finally, we see that if[27] the decays $W \to \widetilde{W}_i \widetilde{Z}_j$ can be excluded at the 1% level, it would be possible to exclude $m_{\widetilde{Z}_3} \leq 70$ GeV.

Future prospects for gaugino searches in Z^0 decays are summarized in Fig. 2, where the minimal visible supersymmetric decay width of the Z^0 is shown as a function of various gaugino masses. For the solid curves, the full set of published constraints has been imposed. Notice that a very light ($m_{\widetilde{Z}_1} \leq 4$ GeV) $\widetilde{Z}_1$ already now implies a sizeable visible supersymmetric branching ratio of the Z-boson; however, the ALEPH bounds[16] are, strictly speaking, not applicable here since it is almost entirely due to $Z^0 \to \widetilde{Z}_2 \widetilde{Z}_3$ decays, which have not yet been searched for. Notice that $m_{\widetilde{Z}_2}$ can at present only be constrained if branching ratios as low as 10^{-5} are detectable, which will (if ever) only be possible with a vastly increased data sample. As seen from Table I, this situation will improve once an improved limit on v/v' or $m_{\widetilde{g}}$ is available.

With such a data sample, other bounds are also likely to improve. In particular, SUSY Higgs bosons with masses up to 50 GeV should eventually be detectable[28] in Z^0 decays at LEP. If we pessimistically assume that none will be found, the bound on v'/v would be sharpened to $v'/v < 0.55$. At the same time, the precision of the measurement of the invisible width of Z is likely to improve,[29] e.g., by "ν-counting" experiments slightly above the peak. We again assume pessimistically that no deviation from the SM will be found, and assume that the bound (1.2) on the invisible width will at best improve to 20 MeV. Under these circumstances, the solid curves change into the long-dashed ones; for the latter, we have ignored the recent ALEPH bound[16] on neutralino production. We see from Fig. 2 that lower bounds of 35 GeV, 70 GeV and 110 GeV on $m_{\widetilde{Z}_2}, m_{\widetilde{Z}_3}$, and $m_{\widetilde{Z}_4}$, respectively, may be attainable from an extension of the ALEPH neutralino search, since it should be possible to identify the unusual decays of $Z \to$ inos at a level $\gtrsim 4 \times 10^{-5}$. We also note that a similar bound on $m_{\widetilde{Z}_3}$ can be obtained if the bound (1.1) is improved to 50 MeV, whereas this does not lead to any limit on $m_{\widetilde{Z}_1}$ or $m_{\widetilde{Z}_2}$ (or the heavy inos).

Note also that the long-dashed $\widetilde{Z}_1$ curve starts at 8 GeV since a lighter $\widetilde{Z}_1$, being Higgsino-like, is excluded because it would result in too large an invisible width. A $\widetilde{Z}_1$ heavier than 8 GeV can be photinolike without violating the $m_{\widetilde{g}} \geq 50$ GeV bound; such a neutralino has only tiny couplings to Z^0 and so would not be excluded by any study of Z^0 decays. We also note that many values of SUSY parameters that led to the degradation[2] of the gluino mass bound obtained[14] by assuming that the gluino decays only via the LSP are now excluded; this, of course, means that the bound on the LSP mass that will emerge from LEP will be even better.[22] (See also Table I.) We stress though that, contrary to a claim in the literature,[30] there is no bound on $m_{\widetilde{Z}_1}$ from collider experiments today, unless the nonobservation of neutralino decays of Z^0 at LEP is explicitly incorporated in the

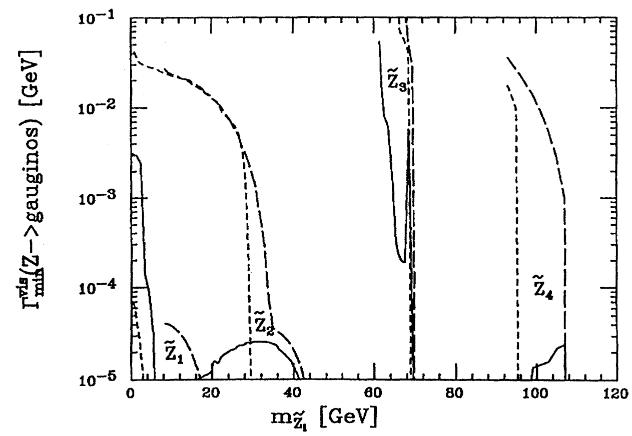

FIG. 2. The minimal total, in principle visible, supersymmetric decay width of the Z, defined as $\Gamma(Z \to \widetilde{W}_- \widetilde{W}_-) + \sum_{i,j} \Gamma(Z \to \widetilde{Z}_i \widetilde{Z}_j)$, as a function of various neutralino masses (the sum includes all neutralino combinations other than $\widetilde{Z}_1 \widetilde{Z}_1$ which is always invisible; note that $\widetilde{Z}_4$ is always too heavy to be produced in real Z^0 decays). The solid curves are valid within the MSSM and for present LEP and $p\bar{p}$ constraints as discussed in the text. The long-dashed curves are also for the MSSM, but with hypothetical future LEP constraints $\delta\Gamma(Z$, nonhadronic$) < 20$ MeV and $v'/v < 0.56$ implemented. The short-dashed curves are also for these sharper constraints, but for the nonminimal model where μ_1/μ_2 is a free parameter. A curve starting at a finite-mass value indicates that smaller values of this mass are incompatible with the imposed constraints. It is instructive to compare the starting point of the long-dashed neutralino curves with the bounds shown in Table I which have been obtained using less restrictive assumptions on the constraints on the gluino mass and the Higgs bosons that may be expected in the future.

analysis.

Finally, the short-dashed curves in Fig. 2 are for the same constraints as the long-dashed ones, but with μ_1/μ_2 as a free parameter rather than as fixed by (2.1). Although this really forms the subject of the next section, we note here that the small $m_{\widetilde{Z}_1}$ region is once again allowed. Notice also, that allowing μ_1 to be free has only a small effect on the bounds on the mass of $\widetilde{Z}_2$ or $\widetilde{Z}_3$.

This concludes our discussion of the MSSM. In the next two sections, we turn to a discussion of how LEP bounds on supersymmetric particles masses are altered in models with further complications. We begin our study with the "free μ_1 model" alluded to in the previous paragraph.

III. GENERALIZATIONS OF THE GAUGINO-HIGGSINO SECTOR

In this section we consider generalizations and extensions of the R-odd spin-$\frac{1}{2}$ sector of the MSSM. In Sec. III A we focus on the simplest modifications which affect only the neutralinos. Since all signals for supersymmetry are, in principle, dependent on the nature of the LSP which can only[8,30,31] be the lightest neutralino if it is also to account for galactic dark matter, even these simple modifications can be of considerable importance;[32] in particular, we will see that the results of the direct searches for neutralinos may be substantially modified. In Sec. III B we construct what, in our opinion, is the minimal extension needed to evade the chargino mass bound obtained from the total width constraint (1.1). As discussed in the Introduction, this is done only with a view to examining the stability of this bound with respect to changes in the model.

A. Models with a generalized neutralino sector

We begin by considering the simplest generalization of the MSSM where the ratios of the SUSY-breaking gaugino masses are not fixed by Eq. (2.1) but are free parameters. Since just the SU(2) gaugino mass μ_2 enters the chargino sector, this has no effect on the chargino bounds other than the fact that the gluino mass constraint (2.4) becomes irrelevant for the analysis. Technically, (2.1) can be evaded[33] by choosing the function[1] $f_{\alpha\beta}$ that determines the kinetic energy terms of gauge superfields to depend nontrivially on the gauge group indices. We stress that there is no good reason to do so (other than the fact that Planck scale physics, which presumably fixes $f_{\alpha\beta}$, is still poorly understood); we study this class of models simply in order to show that this small modification of the MSSM, which only introduces one new parameter into the neutralino sector without the introduction of any new fields, can have major consequences for the regions of parameter space that can be excluded by the published[4,16,20] bounds on chargino and neutralino decays of Z^0.

This is demonstrated in Fig. 3 where the regions excluded by the direct search bounds on $\widetilde{W}_-\widetilde{W}_-$, $\widetilde{Z}_1\widetilde{Z}_2$ and $\widetilde{Z}_2\widetilde{Z}_2$ production cross sections as parametrized in Eqs. (2.10), (2.12), and (2.13) only are plotted for the representative choices $\mu_1/\mu_2 = \frac{1}{2}$ (the MSSM value), $\mu_1/\mu_2 = -\frac{1}{2}$, $\mu_1/\mu_2 = \frac{3}{2}$, and $\mu_1/\mu_2 = \frac{1}{6}$. The region in the upper left corner as well as (except for the $\mu_1/\mu_2 = -\frac{1}{2}$ case) the band between the curves starting at $2m_1 = 0$ and the lowermost curve is excluded by the chargino constraint (2.10), except for the crescent-shaped region, within which $m_{\widetilde{W}_-} - m_{\widetilde{Z}_1}$ is too small or even neg-

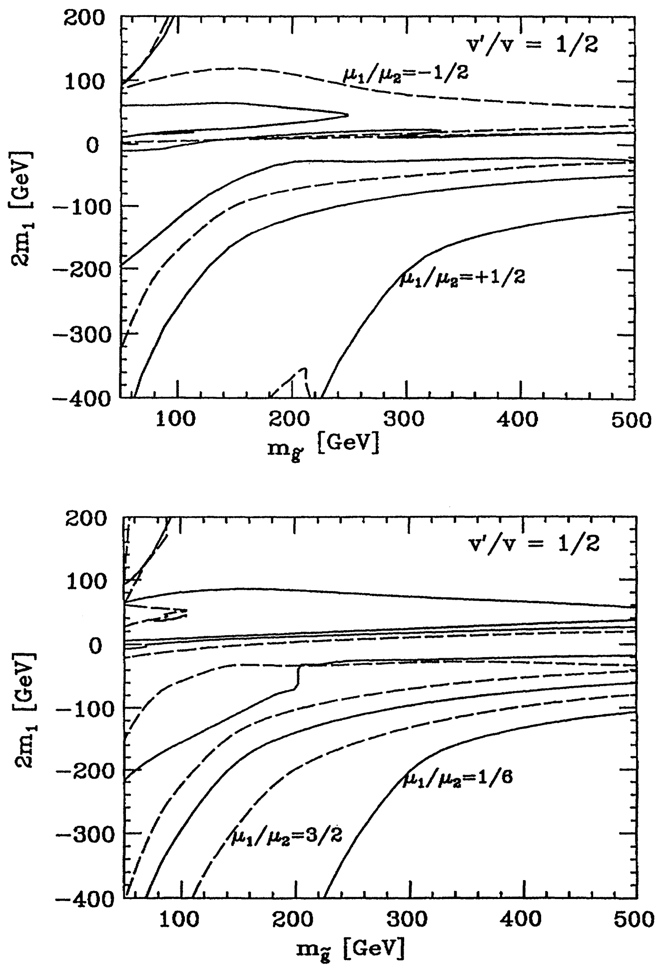

FIG. 3. The regions of parameter space that can be ruled out from the negative results of the direct searches for charginos and neutralinos, as parametrized by Eqs. (2.10) and Eqs. (2.12), (2.13), respectively, for different values of μ_1/μ_2. In Fig. 3(a), $\mu_1/\mu_2 = +\frac{1}{2}$ the MSSM value (solid) and $-\frac{1}{2}$ (dashed), while in Fig. 3(b) $\mu_1/\mu_2 = \frac{1}{6}$ (solid) and $\frac{3}{2}$ (dashed). Note that for $\mu_1/\mu_2 > 0$, most of the narrow allowed region within the large excluded region at negative $2m_1$ is, in fact, excluded by stable charged-particle searches, since in these regions usually $m_{\widetilde{W}_-} < m_{\widetilde{Z}_1}$. Note also that it has always been assumed that the relationship (2.1) between $m_{\tilde{g}} \equiv |\mu_3|$ and μ_2 is unaltered, i.e., $\mu_3 \simeq 3.79\mu_2$.

ative (in which case it is ruled out from stable-particle searches[6] if $\widetilde{W}_-$ is the LSP, or by direct searches[34] if $\widetilde{W}_- \to l\widetilde{\nu}$). For $\mu_1/\mu_2 = -\frac{1}{2}$, this is true everywhere below the lowest dotted hyperbola-shaped region, except for the small triangle at $m_{\tilde{g}} \simeq 200$ GeV, $2m_1 \simeq -400$ GeV. (The base of this triangle increases for even more negative values of $2m_1$.) The nose-shaped region projecting from small $m_{\tilde{g}}$ values near $2m_1 \simeq 50$ GeV is excluded by the neutralino search constraints. Notice that in between this region and the bulk of the "chargino region" discussed above there is a narrow region allowed by the direct-search constraints. For $\mu_1/\mu_2 > 0$, this is due to an accidental cancellation in the $\widetilde{Z}_1\widetilde{Z}_2Z$ coupling whereas, for $\mu_1/\mu_2 = -\frac{1}{2}$, $\widetilde{Z}_1$ and $\widetilde{Z}_2$ are almost degenerate in mass so that the decay products of $\widetilde{Z}_2$ are too soft to give an observable signal. We clearly see that the "excluded regions" are strongly dependent on the value of μ_1/μ_2.

Figure 3 can be understood qualitatively from the observation that for almost all regions of parameter space there is one neutralino with a fairly large ($\geq 1/\sqrt{2}$) U(1)-gaugino component; the mass of this state is determined mostly by μ_1. In all regions that have been probed by direct gaugino searches so far, this neutralino is either $\widetilde{Z}_1$ or $\widetilde{Z}_2$. As long as $\mu_1/\mu_2 > 0$, its mass increases with μ_1. The curves for $\mu_1/\mu_2 = \frac{1}{2}$ therefore lie between the curves for $\mu_1/\mu_2 = \frac{1}{6}$ and $\frac{3}{2}$, respectively. On the other hand, changing the sign of μ_1 decreases the mass of this state if $\mu_2\, 2m_1 v'/v > 0$, and increases it otherwise. Since neutralino searches at present mostly probe the region $2m_1 \geq 0$ (in our convention where $\mu_2, v'/v > 0$), they exclude a much larger region for $\mu_1/\mu_2 = -\frac{1}{2}$ than for the case of the MSSM. On the other hand, in the region of negative $2m_1$ and $m_{\widetilde{W}_-} < M_Z/2$, one often even has $m_{\widetilde{Z}_1} \gtrsim m_{\widetilde{W}_-}$ if $\mu_1/\mu_2 = -\frac{1}{2}$, as mentioned above.

Since the total SUSY width of Z^0 (apart from kinematic effects) is just its width into the "Higgsinos" we may expect that it is less sensitive to the details of the mixing matrix than the direct-search constraints. This is demonstrated in Fig. 4 where we show the region excluded by the constraints (1.1) and (1.2) for the same four choices of μ_1/μ_2. The total width constraint (1.1) contributes most of the excluded region; the constraint (1.2) only excludes the region of very small $2m_1$ and $m_{\tilde{g}} \lesssim 300$ GeV, where $\widetilde{Z}_1$ is Higgsino-like and $m_{\widetilde{W}_-} > M_Z/2$.

Obviously the bound (1.2) becomes effective if the Higgsino-like state becomes lighter than the neutralino with a large U(1) component. This depends on the value of μ_1 and so accounts for the differences in the curves in the small $2m_1$, small $m_{\tilde{g}}$ region. For large values of $m_{\tilde{g}}$, the variation with μ_1/μ_2 is much less as long as the ratio is positive; for $\mu_1/\mu_2 < 0$ and not very small $m_{\tilde{g}}$, a rather large value of $\Gamma(Z \to \widetilde{Z}_1\widetilde{Z}_2)$ always results for small $\mid 2m_1 \mid$ since $\widetilde{Z}_2$ cannot be photinolike as in the case of $\mu_1/\mu_2 > 0$. Thus a somewhat larger region is excluded if $\mu_1/\mu_2 < 0$.

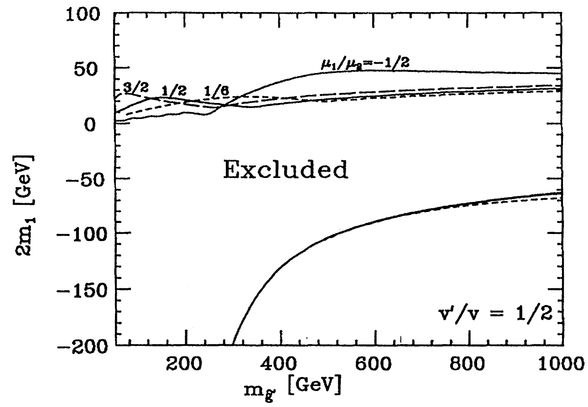

FIG. 4. The region that is excluded by the "inclusive" constraints (1.1) and (1.2), for various values of $\mu_1/\mu_2 = -\frac{1}{2}$ (dotted), $\frac{1}{6}$ (short-dashed), $\frac{1}{2}$ (solid), and $\frac{3}{2}$ (long-dashed). As in Fig. 3, the relation (2.1) between $m_{\tilde{g}}$ and μ_2 is assumed to hold. Note that the region $m_{\widetilde{W}_-} \geq m_{\widetilde{Z}_1}$, which was allowed in Fig. 3, is now excluded for $\mu_1/\mu_2 > 0$.

We see nevertheless that the differences between the various curves in Fig. 4 are much smaller than the corresponding ones in Fig. 3. Note also that, even in these models with μ_1/μ_2 a free parameter, the crescent shaped region in Fig. 3 that was allowed by the direct-search constraints is excluded by the constraints (1.1) and (1.2). This, along with the short-dashed curves in Fig. 2, shows that the region of parameter space as well as the values of $m_{\widetilde{Z}_2}$ and $m_{\widetilde{Z}_3}$ that will be explored at LEP does not differ greatly from that for the MSSM. We emphasize, though, that even in this class of minimally extended models it will be impossible to obtain a bound on $m_{\widetilde{Z}_1}$ if no SUSY signal is ever observed. This is because the additional parameter μ_1 allows us the freedom to make $\widetilde{Z}_1$ a massless U(1) gaugino with all other sparticles heavy by setting $\mu_1 = 0$ and taking both μ_2 and $2m_1$ to be very large.

Yet another way of altering the neutralino sector from that given by the MSSM is to introduce new Higgs supermultiplets, which via the added Higgsino content can alter both the chargino and neutralino mass matrices. As is well known, care must be taken[35] in introducing these fields, since if they couple to both $T_3 = \frac{1}{2}$ and $T_3 = -\frac{1}{2}$ fermions of the SM they can mediate tree-level flavor-changing neutral currents at an unacceptable rate. This is not a problem for an SU(2)×U(1)-singlet superfield n, since it has no gauge-invariant coupling to SM fermions. Such fields are present in models[36] based on the group E_6 and have the virtue that, if they develop a vacuum expectation value (VEV), they can also account for the Higgsino mass $2m_1$ without the addition of a new *ad hoc* mass scale. For these reasons, and primarily because the addition of a singlet is the most economic extension of the gaugino-Higgsino sector, we will consider it in some detail in the following. The effect of the singlet scalars has already been discussed earlier.[37]

The most general purely cubic superpotential for the Higgs sector of the model can be written as

$$f = -ahh'n + \frac{1}{3!}\, bn^3 \qquad (3.1a)$$

so that, effectively, $2m_1 = a < n > \equiv ax$ in this model. The resulting neutralino mass matrix has the form

$$M = \begin{pmatrix} \mu_1 & 0 & \frac{g'v'}{\sqrt{2}} & -\frac{g'v}{\sqrt{2}} & 0 \\ 0 & \mu_2 & \frac{-gv'}{\sqrt{2}} & \frac{gv}{\sqrt{2}} & 0 \\ \frac{g'v'}{\sqrt{2}} & \frac{-gv'}{\sqrt{2}} & 0 & -ax & -av \\ \frac{-g'v}{\sqrt{2}} & \frac{gv}{\sqrt{2}} & -ax & 0 & -av' \\ 0 & 0 & -av & -av' & bx \end{pmatrix} \qquad (3.1b)$$

in the basis $(\lambda_0, \lambda_3, h', h, n)$. It is straightforward to diagonalize M numerically. As might be expected, there is a zero-mass eigenstate if either $\mu_1 = \mu_2 = 0$ (the photino) or $a = 0$ (the doublet Higgsino). It is also possible to obtain a massless eigenstate if the parameters in the mass matrix (3.1b) conspire[32] to make its determinant zero. In general, these states will be complex mixtures of all the gauginos and Higgsinos.

The case $v'/v = 1$ (note that the ALEPH constraint[12] on v'/v no longer applies) lends itself to a particularly simple analysis. In this case, it is easy to check that there is a massless neutralino if either (i) $\mu_2 = \frac{8}{5}g^2v^2/ax$ or, (ii) $bx^2 = -av^2$. Condition (i) has already been obtained in Ref. 32 where it was shown that it describes a massless neutralino with gaugino and doublet-Higgsino components. If $v' \neq v$, this state also picks up a singlet-Higgsino component. If $v'/v = 1$ and condition (ii) is satisfied, the massless neutralino is a pure Higgsino with a singlet component about $av/bx \simeq x/v$ times the doublet components so that its coupling to Z^0 is suppressed by $\sim v^2/x^2$. This state which has no analogue in the MSSM develops small gaugino components if $v'/v \neq 1$.

Notice also that it is possible to have as many as three light neutralinos in this model if $\mu_1 \simeq \mu_2 \simeq ax \simeq bx \simeq 0$. Note also that, just as in the free μ_1 model discussed above, it is easily possible to have a very light LSP that would escape detection at LEP. It is, therefore, worth keeping in mind that any bounds on the LSP mass that may be inferred from the LEP data are special to the MSSM (unless a SUSY signal is actually observed).

The effect of the two extensions of the neutralino sector on the chargino mass bound that can be obtained from the LEP data are shown in Fig. 5 where we have plotted the smallest possible contribution of the charginos and neutralinos to the Z^0 width for (i) the minimal model (dotted), (ii) the "free μ_1/μ_2 model" (dashes) and (iii) the extra singlet model (dot-dashed). Also shown are the contributions from just the charginos with (marked $\widetilde{W}$ only) and without (marked $Z \to \widetilde{h}^+\widetilde{h}^-$) the naturalness constraint $m_{\widetilde{g}} < 1$ TeV. Whenever there are two curves with the same texture, the upper curve is the one obtained by including the direct search limit (2.10) on the chargino mass. The Higgs constraint (2.3) on v'/v has not been included in these curves since it is in general not valid for a non-minimal Higgs-boson sector. The following comments are in order.

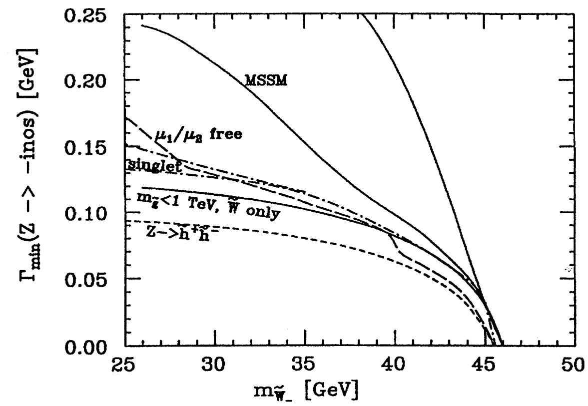

FIG. 5. The minimal total supersymmetric width of the Z boson, $\Gamma(Z \to \widetilde{W}_-\widetilde{W}_-) + \sum_{i,j}\Gamma(Z \to \widetilde{Z}_i\widetilde{Z}_i)$, as a function of $m_{\widetilde{W}_-}$ for various models: MSSM (dotted), the model with free μ_1/μ_2 (long-dashed), and the model with one additional Higgs singlet superfield (dot-dashed). Whenever there are two lines with the same texture, the upper curve has been obtained by imposing the direct-search constraints (2.10), (2.12), and (2.13). For comparison we also show the minimal $\Gamma(Z \to \widetilde{W}_-\widetilde{W}_-)$ for the MSSM, imposing only the constraint $m_{\widetilde{g}} < 1$, TeV (solid curve), as well as $\Gamma(Z \to \widetilde{h}^+\widetilde{h}^-)$, where $\widetilde{h}$ stands for a pure SU(2)-doublet Higgsino.

(i) As already stated, the minimal model bound from (1.1) on the chargino mass with (without) the direct search constraint incorporated is 43 (40) GeV.

(ii) The difference between the minimal model width and that obtained in the two nonminimal models is due to the reduced neutralino decay rate of Z^0 from the additional freedom that is allowed. Nevertheless, we see that even in these models the chargino bound is reduced by only a few GeV.

(iii) The two lower curves obviously give the absolute lowest bound on the chargino mass, regardless of any modifications that affect only the neutralino sector. We see that $m_{\widetilde{W}_-} \gtrsim 37$ GeV if $m_{\widetilde{g}} \lesssim 1$ TeV and that a modest reduction (to about 65 MeV) of the bound (1.1) can lead to absolute bounds exceeding 40 GeV in all models of this class. This should be possible with a Z^0 data sample that exceeds the present one by a factor ≤ 3.

We thus conclude that the chargino limit from LEP is rather robust in that it cannot easily be evaded by simple modifications of the model. As we will see in the next subsection, even more drastic changes leave this bound unchanged. It is to a discussion of this that we now turn our attention.

B. Evading the chargino bound: extended Higgs and gauge sectors

In order to understand how to evade the bound on the chargino mass discussed above, we first study the structure of the $Z\widetilde{W}_-\widetilde{W}_-$ vertex in an SU(2)×U(1) theory with an arbitrary Higgs sector. Since electromagnetic gauge invariance is unbroken, the chargino mass eigenstate is a superposition of the SU(2) gaugino ($\widetilde{\lambda}$) and Higgsino fields (and, if R parity is broken, also lepton fields; it will be evident how this is incorporated) with the same electric charge. In general, the left- and right-handed components $\widetilde{W}_{-L}$ and $\widetilde{W}_{-R}$ will be different superpositions of these fields so we can, in complete generality, write

$$\widetilde{W}_{-L} = U_{11}\ \widetilde{\lambda}_L + \sum_j U_{1j}\widetilde{S}_{jL} + \sum_k U_{1k}\widetilde{D}_{kL} + \sum_l U_{1l}\widetilde{T}_{lL} + \cdots \tag{3.2a}$$

and

$$\widetilde{W}_{-R} = V_{11}\widetilde{\lambda}_R + \sum_j V_{1j}\widetilde{S}_{jR} + \sum_k V_{1k}\widetilde{D}_{kR} + \sum_l V_{1l}\widetilde{T}_{lR} + \cdots . \tag{3.2b}$$

Here U and V are unitary matrices, $j, k, l, \ldots$ run over all the SU(2)-singlet, -doublet, -triplet ,..., Higgsinos with electric charge the same as that of $\widetilde{W}_-$. Since the coupling of Z^0 to a Weyl fermion with weak isospin and charge (T_3, Q) is proportional to $(T_3 - \sin^2\theta_W Q)$, the $Z\widetilde{W}_-\widetilde{W}_-$ coupling is, apart from an overall multiplicative factor of $-g/\cos\theta_W$, given by

$$(-1+\sin^2\theta_W)\left(|U_{11}|^2 \pm |V_{11}|^2\right) + \sum_j \sin^2\theta_W\left(|U_{1j}|^2 \pm |V_{1j}|^2\right) + \sum_k\left(-\frac{1}{2}+\sin^2\theta_W\right)\left(|U_{1k}|^2 \pm |V_{1k}|^2\right) + \sum_l(-1+\sin^2\theta_W)\left(|U_{1l}|^2 \pm |V_{1l}|^2\right) + \cdots, \tag{3.3}$$

where the upper (lower) signs correspond to the vector (axial-vector) coupling. By arranging the mass matrices such that $U = V$ (e.g., by having $v'/v = 1$), it is possible to eliminate the axial-vector coupling (this is also possible by fine-tuning), but the structure of (3.3) clearly demonstrates that for each of the singlet, doublet, triplet (or gaugino) components of the chargino, the mixing factor is always positive. Since $\sin^2\theta_W < \frac{1}{2}$, we conclude that the $Z\widetilde{W}_-\widetilde{W}_-$ vertex can be reduced from its minimum as given by the MSSM only by the introduction of charged SU(2)-singlet Higgsino fields that have substantial mixing with the the gaugino and/or the doublet fields. This charged singlet Higgsino can also, in principle, be the charged gaugino of the right-handed gauge group if it is non-Abelian. Since it is difficult to get these two gauginos to mix in a natural way without mixing the corresponding vector bosons, we will in the following focus on an example where the mixing is due to a charged singlet Higgs field.

In order to obtain such a mixing, we need a $Y = -2$ SU(2)-singlet superfield S interacting with an SU(2)-singlet combination of doublet Higgs fields; since the singlet combination is antisymmetric under interchange of the two doublets, these two doublets must be distinct. Finally, we need a "partner" (S') for S with which it can combine to make a massive Dirac charged field, since a light singlet Higgsino would already have been produced in Z^0 decays via this hypercharge coupling. The simplest superpotential respecting our constraints is

$$f = fS\left(H^+K^0 - H^0K^+\right) + f'S'\left(H'^-K'^0 - K'^-H'^0\right) + \text{mass terms} \tag{3.4}$$

with H and K being $Y = 1$ doublets, S' being a $Y = 2$ singlet and H', K' being $Y = -1$ doublets where $Q = T_3 + \frac{1}{2}Y$. H and H' may be identified with the doublets h and h' of the MSSM. The mixing of the singlet field is induced once the neutral members acquire VEV's due to SU(2)×U(1) breaking. In terms of the $Q = -1$ Dirac fields

$$\begin{aligned} \chi_S &= P_L S - P_R S', \\ \chi_H &= P_L H' - P_R H, \\ \chi_K &= P_L K' - P_R K, \end{aligned} \tag{3.5}$$

the mass matrix in the basis $(\lambda^-,\ \chi_H, \chi_K, \chi_S)$ (where λ^- is the $Q = -1$ gaugino) is

$$M_{\text{chargino}} = \begin{pmatrix} \mu_2 & gv' & gw' & 0 \\ gv & 2m_H & 0 & -fw \\ gw & 0 & 2m_K & fv \\ 0 & -f'w' & f'v' & 2m_S \end{pmatrix}, \tag{3.6}$$

where $2m_H$, and $2m_K$, and $2m_S$ are the HH', KK', and SS' supersymmetric mass terms (analogous to $2m_1$ of the MSSM) and v, w (v', w') are the VEV's of the fields H, K (H', K') and μ_2, as before, is the supersymmetric breaking SU(2) gaugino mass. We note here that we have rotated the fields so that HK' and $H'K$ terms vanish. It is clear that there is enough freedom in (3.6) to allow for a light chargino that has a sufficiently large singlet-Higgsino component so that it does not couple to Z^0.

To see this explicitly, consider $\mu_2 \to \infty$ so that the lighter states are Higgsinos. Then, in the case $f = f', v = v', w = w'$ the state $(0, fw/m_H, -fv/m_K, 1)$ is massless provided $m_S = (fw)^2/m_H + (fv)^2/m_K$ and has zero coupling to Z^0 if $(fw)^2/m_H^2 + (fv)^2/m_K^2 = \sin^2\theta_W/(\frac{1}{2} - \sin^2\theta_W)$.

We stress that we are not advocating this scenario—quite the contrary, in fact. Since considerable modification is required to significantly reduce the $Z\widetilde{W}\widetilde{W}$ coupling from its lowest value as given by the MSSM, we regard the LEP bounds on the chargino mass as more or less model independent. It is nevertheless instructive to note that the field content of our example is (apart from the supersymmetrization) the same as that of a decade old model constructed by Zee to incorporate naturally small off-diagonal Majorana mass terms for the neutrinos.[38] If additional couplings are introduced to incorporate this possibility, a host of lepton-number-violating processes result so that the model can be subjected to experimental scrutiny.

Finally, we observe that a chargino with a vanishing coupling to Z^0 would be produced via its nonresonant electromagnetic couplings with a cross section[3]

$$\sigma = \frac{4}{3}\frac{\pi\alpha^2}{s}\left(1+2\frac{m^2_{\widetilde{W}_-}}{s}\right)\left(1-4\frac{m^2_{\widetilde{W}_-}}{s}\right)^{1/2}$$

which is ~ 10 pb for $m_{\widetilde{W}_-}=0$. Since each of the LEP experiments has already accumulated about 1 pb of integrated luminosity, this electromagnetic component of the cross section should certainly lead to an observable rate in the current LEP run.

IV. EVADING THE SLEPTON BOUNDS FROM Z^0 DECAYS

It is clear that as long as the supersymmetric partners of leptons are also mass eigenstates, their couplings to Z^0 are fixed by electroweak gauge invariance, so that we have an unambiguous prediction for their production rate. If this is indeed the case, then bounds on their masses can only be evaded if their decay modes are very different from those that have been searched for in the LEP experiments. The simplest imaginable possibility is that sleptons decay via $\widetilde{l}\to l\widetilde{Z}_2$, which would be the case if $\widetilde{Z}_1$ is a pure Higgsino [in the MSSM this is only possible if $0.62 < v'/v < 0.77$ because of the invisible width constraint and the constraint (2.3) on v'/v]; then, one may argue that the resulting signature has not been searched for so that the bounds are invalid. It is amply clear that although this may strictly speaking be the case, the resulting event which would contain two hard acollinear leptons accompanied by 2–4 jets $+\not{p}_T$ would never be mistaken for an ordinary event, although it could possibly have been confused with a $t\bar{t}$ event (this would be ruled out since there would be no single lepton events). Similarly, in models where R parity is violated, the LSP decays so that there is no $\not{p}_T$. The signatures for sleptons, although very model dependent, are all spectacular[5], so that we doubt very much that a slepton could have been missed if it had been produced at LEP with the rates as predicted by their SU(2)×U(1) couplings. (Note also that even if it is possible to invent a model where the $\widetilde{l}$ escapes the direct-search experiments, it would still be subject to constraints[8] from the total width and/or peak cross-section measurements.) In our view, the only way a slepton could escape detection in Z^0 decays is if it mixed with an SU(2)-singlet charged field so that it had reduced couplings to the Z^0 as discussed in Sec. III B.

To begin with, let us consider $\widetilde{l}_L$ since there is no bound[8] on $\widetilde{l}_R$ from the width measurements as its coupling to Z^0 is rather small. For definiteness, we will treat only one generation at a time and assume that there is no mixing between generations, as any such mixing is severely constrained. Following our discussion in Sec. III B, we are led to introduce a pair of SU(2)-singlet left-chiral superfields S and S' with $Y(S)=-Y(S')=2$ interacting with the lepton and the $Y=-1$ Higgs superfields via the superpotential

$$f = \lambda_1 h'LS + m_S SS' + \lambda_E h'LE\ , \tag{4.1}$$

where L and E are the SU(2)-doublet and SU(2)-singlet lepton superfields, respectively. Notice that (4.1) allows us to define a conserved lepton number for S and S'. Including soft supersymmetry breaking, trilinear and mass terms gives the mass matrix in the basis $(\widetilde{L}, \widetilde{S}, \widetilde{S}'^*, \widetilde{E}^*)$ as

$$M_{\rm scalar} = \begin{pmatrix} m^2_{\widetilde{L}}+\theta_L^2 m_S^2 & A_1\theta_L m_S & \theta_L m_S^2 & 0 \\ A_1\ \theta_L\ m_S & m^2_{\widetilde{S}}+m_S^2+m_S^2\theta_L^2 & B_S\ m_S & 0 \\ \theta_L m_S^2 & B_S m_S & m^2_{\widetilde{S}'}+m_S^2 & 0 \\ 0 & 0 & 0 & m^2_{\widetilde{E}} \end{pmatrix}, \tag{4.2}$$

where we have neglected the small coupling λ_E, which is responsible for the lepton mass. Here A_1 and $B_S m_S$ are the coefficients of the trilinear and bilinear soft SUSY-breaking terms, $m_{\widetilde{L}}$, $m_{\widetilde{S}}$, $m_{\widetilde{S}'}$ are soft SUSY-breaking masses and

$$\theta_L \equiv \frac{\lambda_1 v'}{m_S}\ . \tag{4.3}$$

It appears from the structure of (4.2) that substantial mixing between $\widetilde{L}$, $\widetilde{S}$, $\widetilde{S}'^*$ is possible while $\widetilde{E}^*$ decouples. Here, special care must be taken since the slepton mixing has been introduced via a superpotential so that the same source of mixing enters the lepton sector via the fermion mass matrix

$$\begin{pmatrix} \lambda_E v' & \lambda_1 v' \\ 0 & m_S \end{pmatrix}, \tag{4.4}$$

where the rows [columns] are in the basis $(L,S')[(E,S)]$. In the limit $\lambda_E\to 0$, the right-handed fermions are unmixed, whereas (for $m_S\gg\lambda_1 v'$) the left-handed fermions have the mixing θ_L given by (4.3), above, independent of the lepton mass.

Mixing of the leptons will alter both their charged- as well as neutral-current couplings. The modifications in the charged-current couplings would show up in the various charged-current processes used to extract the Kobayashi-Maskawa (KM) matrix elements, so that the quoted experimental values of these would include the lepton mixing factor. For example, the range[39] $0.9748 \le | V_{ud} | \le 0.9761$ obtained by a comparison of nuclear β-decay rates and the muon lifetime would be reinterpreted (assuming that θ_L is the only non-SM source of mixing angles) as the range of allowed values of $| V_{ud} | / \cos\theta_{L\mu}$, where $\theta_{L\mu}$ is the mixing of the muon with the singlet fermion. Note that, in this comparison, the electron mixing angle factor $\cos\theta_{Le}$ cancels as it has an identical effect on the muon and nuclear lifetimes. In order to extract a bound on θ_e, we compare the ratio $\Gamma(\pi \to \mu\bar{\nu}_\mu)/\Gamma(\pi \to e\bar{\nu}_e)$ since this depends on $\cos^2\theta_{L\mu}/\cos^2_{\theta_{Le}}$.

We have done an analysis of the mixing angles similar to that in Ref. 40 where we have used $| V_{ub} | < 0.02$ [twice the Particle Data Group[39] (PDG) value] to allow for a possibly large electron mixing angle in the unitarity constraint that gives a lower bound on $| V_{ud} |$. We then find that

$$\theta_{L\mu} \lesssim 0.05, \tag{4.4a}$$

$$\theta_{Le} \lesssim 0.2, \tag{4.4b}$$

whereas the bound $\theta_{L\tau}$ obtained by comparing its experimentally measured lifetime with its theoretically calculated value assuming universality is even weaker. The bound on θ_{Le} is rather mild because, as noted above, it only comes from a measurement of the small branching ratio for the chirally suppressed $\pi \to e\bar{\nu}_e$ decay.

Before turning to the implications of the bounds (4.4) for slepton mixings as given by the mass matrix (4.2), we examine whether the Z^0 coupling to leptons gives any further constraints. It is straightforward to check that the partial width for the decay $Z \to l\bar{l}$ is altered from it SM value by a factor γ given by

$$\gamma = \frac{\frac{1}{4}\cos^4\theta_{Ll} + 2\sin^4\theta_W - \sin^2\theta_W\cos^2\theta_{Ll}}{\frac{1}{4} + 2\sin^4\theta_W - \sin^2\theta_W}, \tag{4.5}$$

where we have neglected the mixing of the right-handed components of the leptons as it vanishes for $m_l = 0$. The LEP experiments[41] have recently measured the separate leptonic branching ratios of Z^0 and found that they are in agreement with SM expectations within a few percent. If we conservatively require $\gamma > 0.95$, we find that

$$\theta_{Ll} \lesssim 0.15 \tag{4.6}$$

which certainly yields the best bound on τ-lepton mixing.

At first glance, the bounds (4.4) and (4.6) suggest that the mixing in the slepton sector is small so that the mass eigenstate of interest can be written as $\frac{1}{N}(1, \epsilon_1, \epsilon_2, 0)$ with

$$\epsilon_1 \simeq \frac{A_1 m_S \theta_L}{m^2_{\tilde{L}} - m^2_{\tilde{S}} - m^2_S}, \qquad \epsilon_2 = \frac{\theta_L m^2_S}{m^2_{\tilde{L}} - m^2_{\tilde{S}'} - m^2_S} \tag{4.7}$$

so that substantial mixing is possible only if the soft SUSY-breaking scalar masses accidently satisfy

$$m^2_{\tilde{L}} \simeq m^2_{\tilde{S}} + m^2_S, \qquad \text{or } m^2_{\tilde{L}} \simeq m^2_{\tilde{S}'} + m^2_S \tag{4.8}$$

which, of course, is fine-tuning at the 10% level. Note also that we must have $m_S \ge 45$ GeV since the SU(2)-singlet charged fermion has not been seen in Z^0 decays. In practice, for substantial slepton mixing, this bound may be considerably higher since, when $\theta_L \ne 0$, Z^0 can decay into a light-heavy fermion pair via a coupling proportional to $\sin 2\theta_L$. This would be signaled by events containing acollinear lepton pairs (with one of the leptons considerably softer than the other) recoiling against a hadron system, an l^+l^- or $l\bar{\nu}$ pair or $\not{p}_T$ since the heavy lepton ψ_l would decay via $\psi_l \to l + Z^*$ or $\nu + W^*$.

Equation (4.8) implies that at least one of $m^2_{\tilde{S}}$, $m^2_{\tilde{S}'}$ must be somewhat negative in order that the lighter slepton state be produced in Z^0 decays. This does not lead to symmetry breaking since the supersymmetric mass m_S more than compensates. Note that the superpotential (4.1) does not contain any large Yukawa couplings of S'—and even the coupling λ_1 is constrained by the lepton mixing to be $\lesssim 10^{-1}$—in order to drive $m^2_{\tilde{S}'}$ to negative values; thus, if as usual we assume there is a common SUSY-breaking scalar mass at the unification scale as in supergravity models, we necessarily must introduce new fields with Yukawa couplings to S' to drive $m^2_{\tilde{S}'}$ to negative values.

Allowing for this, we see that large mixing is, at least in principle, possible. If we denote by $\cos\xi$ the component of $\tilde{l}$ in the mass eigenstate, a calculation similar to that of Sec. III B shows that the coupling of Z^0 to this eigenstate is given by

$$| C | = \frac{g}{\cos\theta_W}\left(\sin^2\theta_W - \frac{1}{2}\cos^2\xi\right) \tag{4.9}$$

so that almost maximal mixing is needed to make C vanish.

As an example, we consider the special case with $A_1 = B_S \simeq 0$. Then, the mixing angle is given by $\tan 2\xi = 2\theta_L m^2_S/(m^2_{\tilde{L}} - m^2_{\tilde{S}'} - m^2_S)$ and the splitting about the mean value $\frac{1}{2}(m^2_{\tilde{L}} + m^2_{\tilde{S}'} + m^2_S)$ is given by $\Delta m^2 = 2\theta_L m^2_S/\sin 2\xi \simeq 2\theta_L m^2_S$ for significant mixing. If we do not allow $m^2_{\tilde{S}'}$ to be arbitrarily negative, we see that although it is possible to get a light scalar that almost does not couple to Z^0, the splitting is just a few tens of GeV so that the associated production of the light-heavy pair is a viable signal once the machine energy is increased to somewhat above $\sqrt{s} = M_Z$.

Note also that the A_1 terms in the mass matrix (4.1) may even be more efficient in mixing the sleptons since A_1 can easily be $2m_S$. These terms cannot be made arbitrarily large since charged breaking minima can develop in the direction $< \widetilde{L}^- >=< \widetilde{S} >=< H' >$ for which the D term vanishes. This can, of course, again be fixed up by introducing new terms in the superpotential which give quartic interactions.

The bounds on the SU(2)-singlet right-handed sleptons can also be evaded by allowing it to mix with sparticles in doublet (or larger) representations of SU(2). As an example, consider the superpotential

$$f = \lambda_E H'LE + \lambda_1 H'DE + m_D DD' \tag{4.10}$$

with D and D' SU(2) doublets with $Y(D) = -Y(D') = -1$. The calculation is very similar to that presented above for the $\widetilde{l}_L$ mixing except, of course, there are no charged-current constraints, so that at least for the mixing of $\widetilde{\mu}_R$ considerably less fine-tuning is required as compared with $\widetilde{\mu}_L$. We will not present any details here.

In summary, we have seen that although it is possible to contrive models where the sleptons mix so as to have vanishing couplings with Z^0, none of these models is particularly believable. All of them involve the introduction of several new fields and some fine-tuning. It is amusing that the class of models we have constructed to evade the slepton bounds from LEP all predict a new scalar that should rule them out as soon as the center-of-mass energy is increased by a few tens of GeV. In conclusion, it appears to us that the bounds on the charged-slepton masses from the LEP experiment are very difficult to evade except in very contrived scenarios.

V. SUMMARY AND OUTLOOK

Last year, the four experiments at LEP together accumulated $\sim 10^5$ Z^0 decays. Detailed analyses of this data sample have shown[9] remarkable agreement with the SM. As more data is accumulated, various properties of the Z^0 boson (mass and partial width for decays into different channels) will be measured with even higher precision. Comparison of theory and experiment will then entail[42] a knowledge of radiative corrections which, of course, involve the properties of all the particles that couple to Z^0, whether or not they are kinematically accessible at LEP energy. One way to test the SM in the near future would be to ask whether the values of all measurements are consistent with a unique set of values for the t-quark (m_t) and Higgs-boson (M_H) masses, which are the two unknown parameters of the SM. It is conceivable that two different measurements may require incompatible values for (m_t, M_H), in which case we would conclude that the experiment is in conflict with the SM.

An alternative strategy would be to look directly for new particles that are produced in Z^0 decays. This has been the main strategy adopted by the LEP experiments at least up to now. The reason, of course, is that these analyses are more straightforward and do not require as large a data sample. The obvious disadvantage is that the domain of masses that can be searched for is limited by M_Z and, in the many cases where the new particles can only be pair produced, by $M_Z/2$.

The results[4,16] of these searches have been used to put lower bounds on masses of various particles including new quarks and leptons, the Higgs boson and supersymmetric particles. These bounds depend on the strategy used to search for these particles which, in turn, depends on how the particles decay. Thus, to a lesser or greater extent, such bounds are all model dependent.

This is especially true for supersymmetry searches because a heavy sparticle typically has several decay modes possible, each one involving a supersymmetric daughter, if we assume as usual that R parity is conserved. The daughter then decays and the chain ends in the stable LSP. Clearly the signals for the production of any sparticle then depend on its couplings to Z^0 as well as the properties of other sparticles lighter than itself. A major purpose of this paper is to study just how the various lower bounds on sparticle masses resulting from the negative searches at the LEP collider depend on the assumptions of the MSSM, which is the framework usually adopted in such studies. Toward this end we have attempted to construct models to circumvent these bounds. If this can be done by a minor change of the minimal model, then the bound in question is obviously very model dependent, whereas if the modifications needed are many and contrived, we will regard the bound as robust. Needless to say, we do not advocate any of the models we have constructed for this purpose.

Throughout, we have assumed that R-parity is conserved since sparticle production in R-parity-violating models is known to lead to events with spectacular collider signatures involving multiple leptons and/or jets. Although such events have not been specially searched for at LEP, we believe that they would have been sufficiently striking so as not to escape notice in these experiments.

At this point, it is worth reiterating that lower bounds on sparticle masses (and the mass of anything accessible in Z^0 decays) can also be obtained[8] from the LEP measurement of the total width and peak hadronic cross section as new decays tend to increase the former and reduce the latter. Although these bounds are not as strong as the bounds from direct searches, they are independent of how the new particles decay (the bounds from the peak cross section require, however, that hadronic modes do not dominate the decays of the new particle).

Section III deals with the bounds on supersymmetric particles within the framework of the minimal model. After a brief discussion of the theoretical framework, we have proceeded to a critique of bounds from various LEP data. In the process, we have upgraded our earlier analysis[8] on the bounds from the inclusive decays of Z^0. More importantly, we have illustrated how these inclusive bounds can be used in conjunction with bounds from direct searches to obtain a stronger bound than would have been possible on the basis of any one strategy alone. The best example of that is the improvement of the MSSM chargino bound from 40 GeV (2.11a) to 43 GeV (2.11b) regardless of the decay patterns of the chargino, whereas the exclusive bound had been obtained assuming that the chargino could only decay via $\widetilde{W} \to f\bar{f}\widetilde{Z}_1$ to the LSP. The other bounds are $m_{\widetilde{\nu}} > 28.9(38.4)$ GeV for one (three degenerate) light sneutrinos, which, when taken together with the ALEPH constraint[12] $v'/v < 0.77$, implies that $m_{\widetilde{l}_L} > 49.8(55.8)$ GeV, but only within the framework of the MSSM. The strongest bound[4] on $\widetilde{l}_R$ comes from direct searches assuming $\widetilde{l}_R \to l\widetilde{Z}_1$ is dominant. We have seen that there are regions of parameter space of the MSSM where $\widetilde{l}_R \to l\widetilde{Z}_2$ essentially all the time; nevertheless, we have argued that the events would be very distinctive at least in the case of $\widetilde{e}_R$ and $\widetilde{\mu}_R$, so it is unlikely that these sparticles would have escaped detection. We have also argued that $\widetilde{\tau}_R\widetilde{\tau}_R$ pairs, which would lead to high sphericity, high-multiplicity events with abnormally narrow jets (from τ misidentification), would (even though not as distinctive as $\widetilde{e}_R$ or $\widetilde{\mu}_R$ events) be readily identifiable, so that a lower limit of about 40 GeV on $\widetilde{l}_R$ masses is probably valid. We have also seen that even within the MSSM a bound on $m_{\widetilde{Z}_1}$ can be derived from the published data only if the decays $Z \to \widetilde{Z}_1\widetilde{Z}_3$ are incorporated into the analysis and it is further assumed that $\widetilde{Z}_3$ decays directly to the LSP. The bound on $m_{\widetilde{Z}_2}$ is simply governed by that on $m_{\widetilde{g}}$ and is given by $m_{\widetilde{Z}_2} \geq m_{\widetilde{g}}/6$. Finally, $m_{\widetilde{Z}_3}$ is constrained to be larger than 61 GeV.

What about the future? Clearly, the most important task would be to constrain (or measure) the $\widetilde{Z}_1$ and $\widetilde{Z}_2$ masses. We can see from Fig. 2 that either (i) the Higgs boson will be discovered, (ii) the invisible width of Z^0 will be too large (this can be attributed to $Z \to \widetilde{Z}_1\widetilde{Z}_1$), or (iii) lower bounds of 8 GeV and 35 GeV would emerge on $m_{\widetilde{Z}_1}$ and $m_{\widetilde{Z}_2}$. As can be seen from Table I, the former bound would be significantly improved if the hadron collider experiments can improve the bound $m_{\widetilde{g}} > 50$ GeV, independently of how the gluino decays (note that this should be possible[2] since very light charginos are now ruled out). Any lower bound on $m_{\widetilde{Z}_1}$ would obviously have cosmological significance. Finally, we have also seen that if we are lucky (e.g., if $m_{\widetilde{\nu}} \sim M_Z/2$), a study of W boson decays at the Fermilab Tevatron would lead to even stronger constraints on SUSY parameters. Notice that the most stringent limits emerge by putting together various constraints. It is here that we see supersymmetry at work, for instance, in the observation that if the Higgs boson is too heavy, there must not be a light Higgsino.

Section III deals with the evasion of the bounds on chargino and neutralino masses in nonminimal models. We have seen that the neutralino couplings (and hence any potential mass bounds on them) are very sensitive to even small changes in the assumptions of the minimal model. The two simple alterations of the neutralino sector that we have considered are (i) giving up[33] the assumption that the SU(2) and U(1) gaugino masses are related by Eq. (2.1), the unification condition usually assumed, and (ii) introducing one extra SU(3)×SU(2)×U(1) singlet superfield. Such fields are often present in models with large gauge groups.[36]

Neither of these alter the chargino sector of the theory. The effect of relaxing either of these is seen in Figs 3–5. Figure 3 shows that altering μ_1/μ_2 from its value in the MSSM drastically changes the regions of SUSY parameter space ($m_{\widetilde{g}}$ is related to μ_2) that can be excluded by the experiments which directly search for the neutralino decays of Z^0, whereas from Fig. 4 we see that the total variation in the inclusive search is much smaller. From Fig. 5, it is clear from the difference in the curves for the minimal model and those for the free μ_1 case or for the model with the singlet that the neutralino contribution to the Z width (recall that the whole difference is due to neutralinos as the chargino sector is untouched) can be drastically reduced. We thus conclude that any bounds on neutralino parameters obtained by the LEP experiments are strongly dependent on the assumptions of the MSSM. Also, in these extended models, the LSP could be light and provide the dark matter and yet escape detection at colliders.

Charginos in nonminimal models have been studied in Sec. III B where we have shown that in any theory with an (effective) SU(2)×U(1) symmetry, the LEP bounds can be evaded only if the chargino contains a substantial component of an SU(2) singlet charged field that mixes with the Higgsino or gaugino fields. We saw that in order to bring about this mixing we were forced to introduce a pair of SU(2) singlets (S, S') and another pair (K, K') of SU(2)-doublet superfields. We then saw that in order to get a light chargino that did not couple to Z^0, the parameters had to be adjusted to satisfy at least one *ad hoc* relation. We were, therefore, led to conclude that the bound on $m_{\widetilde{W}_-}$ obtained from the LEP experiments assuming that the chargino is a doublet Higgsino is valid, except in a class of very contrived models.

Finally, in Sec. IV we considered whether it is at all possible to avoid the bounds on slepton masses. The same arguments as for charginos imply that this is possible only if $\widetilde{l}_L$ mix with a charged singlet and $\widetilde{l}_R$ with any higher representation of $SU(2)_L$. As we saw, the situation here is further complicated by the fact that in a supersymmetric theory the slepton and lepton mixings are proportional. The latter are, of course, constrained by the observed lepton universality. Somewhat surprisingly, we found that this constraint was really stringent only for muons but that, for e and τ, mixing angles ~ 0.15 are quite acceptable. In order to realize the mixing we were led to introduce a pair of charged SU(2)-singlet superfields (S, S') and forced to fine-tune at least one of the SUSY-breaking masses to $\sim 10\%$ (more for smuons) in order for the slepton mixing to be substantial without violating constraints on lepton mixing. Further, we saw that new Yukawa couplings of S and S' were needed, which meant introducing even more superfields. This was needed to get the SUSY-breaking scalar $m^2_{s'}$ to be negative, which in turn was necessary for the new singlet SUSY fermion to be heavy enough to avoid detection while the slepton could still be light. The situation for right-handed sleptons was very similar. Our conclusion was that these models, if anything, were even uglier[43] than the models contrived to escape the chargino bound.

To conclude, the direct searches for SUSY particles at LEP have led to bounds close to the kinematic limit on the masses of charged sparticles. Within the framework of the MSSM, these bounds can be strengthened if taken in conjunction with the inclusive bounds from the constraints (1.1) and (1.2). These constraints also lead to a bound on the sneutrino mass. We have also seen that while simple changes in the model can cause the bounds[16] on neutralino decays of Z^0 to disappear, it is only in very ugly and contrived models that the LEP bounds on the chargino or slepton mass are substantially weakened. Whether or not one finds new physics either by direct searches for Z^0 decays (with small branching fractions) or by finding deviations from the SM in precision measurements remains to be seen.

Note added in proof. Since this paper was submitted, the CDF Collaboration has announced[44] a preliminary lower limit of about 150 GeV on the gluino mass under the assumption that the gluino only decays into the USP. It has been estimated[45] that the incorporation of the non-LSP decays of the gluino as given by the MSSM reduce this bound to about 135 GeV. We see from Table I that this implies a bound of about 15 GeV on the LSP mass.

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