Variations on minimal gauge-mediated supersymmetry breaking

Michael Dine
Santa Cruz Institute for Particle Physics, University of California, Santa Cruz, California 95064

Yosef Nir
Department of Particle Physics, Weizmann Institute of Science, Rehovot 76100, Israel

Yuri Shirman
Santa Cruz Institute for Particle Physics, University of California, Santa Cruz, California 95064

(Received 14 August 1996)

We study various modifications to the minimal models of gauge-mediated supersymmetry breaking. We argue that, under reasonable assumptions, the structure of the messenger sector is rather restricted. We investigate the effects of possible mixing between messenger and ordinary squark and slepton fields and, in particular, violation of universality. We show that acceptable values for the μ and B parameters can naturally arise from discrete, possibly horizontal, symmetries. We claim that in models where the supersymmetry-breaking parameters A and B vanish at the tree level, $\tan\beta$ could be large without fine-tuning. We explain how the supersymmetric CP problem is solved in such models. [S0556-2821(97)03603-5]

PACS number(s): 11.30.Pb, 11.15.Ex, 11.30.Er

I. INTRODUCTION

Most speculation about supersymmetry phenomenology starts with the assumption that supersymmetry is broken at an extremely large energy scale, of order 10^{11} GeV, and that the breaking is fed down to the partners of ordinary fields through gravitational interactions. There has been renewed interest, recently, in the possibility that supersymmetry might be broken at much lower energies, of the order of 10's to 100's of TeV. This interest has grown out of an appreciation of the supersymmetric flavor problem, as well as out of successful efforts to build models with dynamical supersymmetry breaking at low energies [1–2]. More recently, it has also been fueled by one small piece of experimental support: a single $e^+e^-\gamma\gamma\not{E}_T$ event observed at the Collider Detector at Fermilob (CDF) [3–8].

Existing models of low energy supersymmetry breaking assume that gauge interactions are the messengers of supersymmetry breaking. This mechanism is referred to as ''gauge-mediated supersymmetry breaking'' (GMSB). Such models are highly predictive. Indeed, all 106 new parameters of the minimal supersymmetric standard model are typically predicted in terms of two or three new parameters. For example, the simplest model [so-called ''minimal gauge mediation'' (MGM)] possesses a messenger sector consisting of a single $\mathbf{5}+\underline{\mathbf{5}}$ of SU(5), i.e., color triplets, $q+\overline{q}$, and weak doublets, $l+l$. These couple to a single gauge-singlet field S through a superpotential

$$W=\lambda_1 Sq\overline{q}+\lambda_2 Sl\overline{l}. \tag{1.1}$$

The field S has a nonzero expectation value both for its scalar and auxiliary components S and F_S. Integrating out the messenger sector gives rise to gaugino masses at one loop and scalar masses at two loops. For the gauginos, one has

$$m_{\lambda_i}=c_i\frac{\alpha_i}{4\pi}\Lambda, \tag{1.2}$$

where $\Lambda=F_S/S$, $c_1=5/3$, $c_2=c_3=1$, and $\alpha_1=\alpha/\cos^2\theta_W$. For the scalar masses one has

$$\widetilde{m}^2=2\Lambda^2\left[C_3\left(\frac{\alpha_3}{4\pi}\right)^2+C_2\left(\frac{\alpha_2}{4\pi}\right)^2+\frac{5}{3}\left(\frac{Y}{2}\right)^2\left(\frac{\alpha_1}{4\pi}\right)^2\right], \tag{1.3}$$

where $C_3=4/3$ for color triplets and zero for singlets, $C_2=3/4$ for weak doublets and zero for singlets, and $Y=2(Q-T_3)$ is the ordinary hypercharge.[1]

Because the scalar masses are functions of only gauge quantum numbers, these models also automatically solve the supersymmetric flavor problem. This feature is preserved in any theory in which gauge interactions are the messengers of supersymmetry breaking. As a result, such models do not suffer from flavor-changing neutral currents and can naturally have small CP violation.

One can argue, based on these features alone, that low energy supersymmetry breaking is in many ways more appealing than models with intermediate scale breaking. In fact, it is fair to say there do not yet exist computable models of intermediate scale breaking.[2] On the other hand, while

[1]These formulas predict a near degeneracy of the B-ino and the right-handed sleptons. Important corrections due to operator renormalization and D terms have been discussed in [9].

[2]Supergravity models are nonrenormalizable, and so without some underlying finite theory, none of the soft breakings can be determined, even if the supersymmetry-breaking mechanism is specified. In string theory, one cannot compute the soft breaking terms without understanding the dynamics which selects a particular point in the moduli space. While regions of the moduli space in which string theory might yield squark degeneracy have been identified [10–12], it is difficult to understand why these regions would be preferred.

successful models of low energy breaking have been constructed, it would be difficult to claim that we are yet in possession of the analogue of the Weinberg-Salam model for supersymmetry, a model compelling for its elegance and simplicity.

The mass formulas of Eqs. (1.2) and (1.3) are remarkably predictive. But given that we do not yet possess a compelling model, it is natural to ask in what sense such formulas are inevitable consequences of low energy supersymmetry breaking. In Ref. [4], some plausible modifications of these formulas were mentioned. In this paper, we will attempt a more systematic analysis of this issue. In Sec. II, we will consider weakly coupled models. In such theories, some very modest assumptions severely restrict the allowed possibilities. Gauge mediation must play a dominant role, and the messenger sector must consist of small numbers of vector-like representations of SU(5). As already discussed in Ref. [4], the overall coefficients in Eqs. (1.3) and (1.2) may change.

But we also find that it is possible to obtain departures from universality.[3] In particular, in models such as those of Ref. [2], it has been assumed that the messenger sector is completely separated from the visible sector. This can be assured by discrete symmetries. However, one can consider relaxing this condition. Indeed, one might well want to since otherwise the models possess stable particles which are problematic in cosmology.[4] If one allows mixing, there are additional, nonuniversal contributions to scalar masses. We will evaluate these contributions in Sec. III and find that they are negative and proportional to the squares of a new set of Yukawa couplings. One might worry that, as a result, they will spoil the good features of gauge mediation. However, with a minimal messenger sector, i.e., one pair of either $5+\bar{5}$ or $10+\overline{10}$, only masses of one sfermion generation are shifted and the first two generations are likely to remain degenerate. Furthermore, we might expect that, similarly to the ordinary quark and lepton Yukawa coupling matrix, many of these couplings are small, and so only a few states will show departures from universality.

In the MGM model of Ref. [2], all supersymmetry-breaking scalar and gaugino masses depend on one parameter only. The generation of a $\mu H_U H_D$ term in the superpotential and the generation of a supersymmetry-breaking $BH_U H_D$ term in the scalar potential require independent mechanisms. Furthermore, the mechanism presented in [2] for generating B involves fine-tuning of order $(\alpha_2/\pi)^2$. It was suggested that a discrete (possibly horizontal) symmetry could account for the magnitude of B and μ, but no concrete model was presented. In Sec. IV we examine the question of naturalness in more detail. We present a specific version of the minimal model of Ref. [9] where a discrete symmetry predicts μ and B terms of the correct order of magnitude. Previous, related studies, were made in [14,15].

One of the surprising results of this analysis is the fact that a large $\tan\beta$ arises naturally. In Refs. [16–18] dimensional analysis was used to argue that a large $\tan\beta$ requires (in models with two Higgs doublets) fine-tuning of order $1/\tan\beta$ in order to avoid unacceptably light charginos. In Sec. V, we point out that the existence of several energy scales in the full high energy theory can invalidate this analysis and, in particular, that it need not hold in models of low energy supersymmetry breaking.

Another nice feature of the minimal messenger model of Ref. [9], where A and B vanish at the tree level, is that the supersymmetric CP-violating phases ϕ_A and ϕ_B vanish. The supersymmetric CP problem, namely, the $\sim 10^{-2}$ fine-tuning required in generic supersymmetric models to satisfy constraints from electric dipole moments, is then solved. We briefly discuss this point in Sec. VI.

It is quite possible that the dynamics which breaks supersymmetry is strongly coupled. This is an area which has only been partially explored [4]. For such theories, it is more difficult to list general constraints. We will not make a serious effort to tackle this problem here, but we will at least enumerate some of the issues in our concluding section, Sec. VII.

II. CONSTRAINTS ON THE MESSENGER SECTORS OF WEAKLY COUPLED MODELS

It is possible to construct models of low energy supersymmetry breaking using the O'Raifeartaigh and/or Fayet-Iliopoulos mechanisms, in which all couplings are weak and which can be analyzed in perturbation theory. One can imagine that the required couplings are small parameters generated by some more microscopic theory. This microscopic theory might be of the type discussed in Ref. [2], in which dynamical supersymmetry breaking at a not too distant scale generates such terms in an effective action for the messengers. Or one could imagine that it is a theory such as string theory and that the small mass scale is generated by tiny nonperturbative string effects.

In this section, we will not worry about the detailed origin of these terms, but instead ask about the phenomenological constraints on the messenger sector. In a theory which is weakly coupled,[5] it is possible to prove a number of general results. Dimopoulos and Georgi [19] showed long ago that, as a consequence of sum rules, one cannot obtain a realistic spectrum at the tree level in any globally supersymmetric theory. This means that at least some masses must be generated radiatively. In such a theory, some set of fields, which we will call the ''messengers,'' must feel the breaking of supersymmetry at the tree level. Ordinary fields will couple to these. One might imagine that the messengers could all be

[3]In this paper, we will use the term ''universality'' to mean scalar masses which are functions only of gauge quantum number and A terms which are small or proportional to fermion Yukawa couplings.

[4]In [13] it is shown that under certain circumstances, these particles are suitable dark matter candidates. However, the lightest of these needs to be quite light, of order 5 TeV (compared to a natural scale of 30 TeV or more). This potentially represents a fine-tuning of 1 part in 30 or worse. As these authors note, other dark matter candidates are likely to be found elsewhere in these models.

[5]This does not necessarily mean that supersymmetry breaking arises in perturbation theory. Models in which the hidden sector dynamics is calculable semiclassically, for example, would fall in this class.

neutral under ordinary gauge interactions, but it is easy to rule out this possibility. This is because only Higgs fields have the correct quantum numbers to couple (through renormalizable interactions) to the messengers.[6] But this means that Higgs boson masses squared will arise at lower order (by several loops) than gaugino masses, and so gluino masses will be far too small.

So we see that the messenger sector must contain fields which are charged under the standard model group. These fields must come in vectorlike representations. This simply follows from the fact that these masses must be much larger than the weak scale. If we require perturbative coupling unification with a desert, they must come in complete SU(5) multiplets. (For a different scenario, see Ref. [20].) Moreover, we can require that the couplings remain perturbative at least up to the grand unified theory (GUT) scale. This means that one can allow at most four $\mathbf{5}+\overline{\mathbf{5}}$'s or one $\mathbf{10}+\overline{\mathbf{10}}$. SU(5) adjoints are not allowed.

Next, we must ask to what the messengers can couple. In order that they obtain large masses, the messengers almost certainly must couple to fields in the superpotential which obtain vacuum expectation values (VEV's).[7] These fields must be gauge singlets. The simplest possibility, as in the models of Ref. [2], is that the F components of these fields also have expectation values. These F components might also arise at the tree level or through loop corrections (e.g., mixing terms in the Kahler potential; see, for example, Ref. [21]). This will lead to formulas which are simple modifications, depending on the number of messenger fields, of Eqs. (1.2) and (1.3). Alternatively, the singlets might have vanishing F components. The messengers might acquire supersymmetry-breaking masses through loops, either involving superpotential couplings or gauge interactions. Aesthetic issues aside, such models will have difficulty explaining the $\gamma\gamma$ events, should they turn out to be real, since the scale of Goldstino decay constant will tend to be rather large and the (next to lightest supersymmetric particle NLSP) will tend to decay outside the detector. (This is also an issue in models in which the F component of the scalar field arises in loops.) Such models will still lead to mass formulas somewhat different in form than those of Eqs. (1.2) and (1.3). (Such models appeared in Ref. [1].) Of course, masses are still functions only of gauge quantum numbers. In this case, the number of soft breaking parameters is equal to 8, plus the μ and B terms.

So it seems most likely that the messenger sector will consist of some number N_5 of $\mathbf{5}+\overline{\mathbf{5}}$ representations ($N_5<5$) or one $\mathbf{10}+\overline{\mathbf{10}}$ representation, coupling to some number N_S of singlet fields with nonvanishing scalar and F components. If there are N_5 $\mathbf{5}+\overline{\mathbf{5}}$'s and N_S singlets, and we assume that there is no mixing of the messenger fields with ordinary fields, the superpotential in the hidden sector has the form

$$\sum_{i=1}^{N_S}\sum_{J,K=1}^{N_5}(\lambda_1)_{iJK}S^i\overline{q_I}q_J+(\lambda_2)_{iJK}S^i\overline{l_I}l_K. \tag{2.1}$$

For large enough N_S and N_5, the masses of squarks, sleptons, and gauginos, with given gauge quantum numbers, become independent parameters. We might still expect that their masses would be arranged hierarchically as in Eqs. (1.2) and (1.3), but this is not necessarily the case. For example, take $N_S=2$ and $N_5=1$. Suppose that S_1 has a large scalar component and a small F component, while S_2 has a small scalar component and a large F component. Couplings to S_1 could give all messenger quarks and leptons comparable supersymmetry-conserving masses, while couplings to S_2 could be, say, order 1 for messenger quarks, but small for messenger leptons. This could alter the hierarchy (in a universal way), giving a much larger than expected ratio of squark to slepton masses. Similarly, one could arrange that doublets are lighter than singlets or that the gaugino hierarchy is altered.

On the other hand, the modifications of the hierarchy cannot be too drastic or one will face other problems. Given the experimental constraints on squark masses, one cannot take squarks much lighter than lepton doublets, without having to fine-tune Higgs parameters. Similarly, if squarks are extremely heavy, one will have an extremely large, negative contribution to the Higgs boson masses and further problems with fine-tuning. Finally, if one wants to explain the $\gamma\gamma\not{E}_T$ events in this framework, the fundamental scale of supersymmetry breaking cannot be much larger than 10^3–10^4 TeV. Still, it is worth keeping in mind that the hierarchy of squark and gaugino masses suggested by the MGM need not hold, even in weakly coupled theories, provided that they are sufficiently complicated. It is a simple matter to perform the analogous analysis when the messenger sector contains a single $\mathbf{10}+\overline{\mathbf{10}}$.

So far, we have explained how modifications of the hierarchy might arise, but not violations of universality. The fields $\overline{q}$ have the same quantum numbers as the ordinary d fields. We define d as the three fields that do not have a couplings of the form (1.1) or (2.1). In the models of Ref. [2] (and most other recent works), it was implicitly assumed that only the d fields have Yukawa couplings, H_DQd. Indeed, terms of the form $H_DQ\overline{q}$ can be forbidden by discrete symmetries. (Analogous comments hold for the lepton fields or for $\mathbf{10}+\overline{\mathbf{10}}$ messengers.) On the other hand, such Yukawa couplings may be present and can lead to more profound modifications of the minimal-gauge-mediated theory than we have contemplated up to now. Moreover, in the absence of these couplings, the messenger sector contains stable or nearly stable particles, which may be problematic in cosmology. In the next section, we explore the consequences of introducing such couplings.

III. MESSENGER-MATTER MIXING

In order to understand possible modifications of the spectrum in the presence of mixing between messenger fields and ordinary matter fields, consider first the model of Eq. (1.1). The VEV $\langle S\rangle$ gives a supersymmetric contribution to the mass of the messenger quarks and leptons, while $\langle F_S\rangle$ leads to a supersymmetry-violating splitting in these multiplets. At

[6]We are assuming here that R parity is conserved, but all of the remarks which follow are easily modified in the case of broken R parity.

[7]Alternatively, there might be "bare masses" in the superpotential analogous to the μ term. These might arise by the mechanism for the μ term described in Ref. [2] and, further, in Sec. IV.

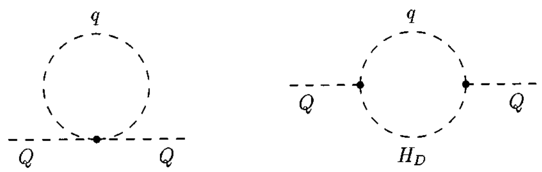

FIG. 1. Scalar-loop contributions to squark mass shifts. Q are ordinary left-handed squark doublets, H_D is the down Higgs doublet, and q are the messenger squarks.

one loop, gauginos gain mass through their couplings to these fields; at two loops, ordinary squarks and sleptons gain mass. In Eqs. (1.2) and (1.3), the parameter Λ is given by $\Lambda = F_S/S$ (here and below, expectation values are understood).

The simple modification that we consider takes place in the Yukawa sector. The messenger ℓ field has the same gauge quantum numbers as the ordinary lepton doublets; $\bar{q}$ has the same quantum numbers as the $\bar{d}$ quarks. Thus, in the absence of a symmetry, one expects these fields to mix. In particular, in the Yukawa couplings

$$H_D L_i Y^{\ell}_{ij} \bar{e}_j + H_D Q_i Y^{d}_{ij} \bar{d}_j, \tag{3.1}$$

each of L_i and $\bar{d}_i$ refers to the four objects with the same quantum numbers. Then Y^{ℓ} is a 4×3 matrix while Y^d is a 3×4 matrix. By convention, we call L_4 and d_4 the linear combination of fields which couple to S in Eq. (1.1). We refer to Y^{ℓ}_{4i} and Y^{d}_{i4} as exotic Yukawa couplings.

The exotic Yukawa couplings contribute, through one-loop diagrams, to the masses of the ordinary squarks and sleptons. These diagrams are indicated in Fig. 1.

It is a simple matter to compute these in a power series in F_S/S^2. The zeroth order term, of course, vanishes by supersymmetry. The first order term vanishes as a result of an accidental cancellation. In order to understand the result, suppose first that only one of the sleptons, say, $\bar{e}_3$, has a substantial Yukawa coupling to L_4 and call this coupling y_{ℓ}. [There is actually no loss of generality here. In general, the affected slepton is the combination $\Sigma_i Y^{\ell}_{4i} \bar{e}_i$, with $y^2_{\ell} = \Sigma_i (Y^{\ell}_{4i})^2$.] The mass shift is

$$\delta m^2_{\bar{e}_3} = -\frac{M^2}{6}\frac{y^2_{\ell}}{16\pi^2}\frac{|F_S|^4}{M^8}, \tag{3.2}$$

where $M=\lambda_2 S$ is the mean mass of the ℓ_4 multiplet. Using Eq. (1.3) for the (universal) two-loop contribution to $m^2_{\bar{e}}$, we find

$$\frac{\delta m^2_{\bar{e}_3}}{m^2_{\bar{e}}} = -\frac{1}{12}\frac{y^2_{\ell}}{\alpha^2_Y}\frac{|F_S|^2}{M^4} \approx -10^3 y^2_{\ell}\frac{|F_S|^2}{M^4}. \tag{3.3}$$

There is also a related shift of the down Higgs boson mass:

$$\frac{\delta m^2_{H_D}}{m^2_{H_D}} = -\frac{1}{9}\frac{y^2_{\ell}}{\alpha^2_2}\frac{|F_S|^2}{M^4} \approx -10^2 y^2_{\ell}\frac{|F_S|^2}{M^4}. \tag{3.4}$$

A few comments are in order, regarding the results (3.3) and (3.4).

(i) Since the result of Eq. (3.2) is proportional to $|F_S|^4$, in contrast to the two-loop contribution of Eq. (1.3) which is proportional to $|F_S|^2$, there is a natural way of understanding how, even for Yukawa couplings of order 1, one-loop corrections could be comparable to two-loop gauge corrections, rather than much larger. Explicitly, $\delta m^2_{\bar{e}_3}/m^2_{\bar{e}} \lesssim 1$ if $|F_S|/M^2 \lesssim 0.03$.

(ii) Related to (i), it is important that contributions to masses of squarks and sleptons not be too large, or charged or colored fields will obtain expectation values. For our example above, the negative correction to the Higgs boson mass is of the same order. But given that the correction to the singlet cannot be too large, the fractional correction to the doublet mass will be rather small.

(iii) With a single messenger $\mathbf{5}+\bar{\mathbf{5}}$ pair, this mass shift affects only one right-handed slepton generation. The other two remain degenerate.

(iv) If, similarly to ordinary Yukawa couplings, $Y^{\ell}_{4\tau} \gg Y^{\ell}_{4i}$ for $i=e,\mu$, then the shift is in the mass of the right-handed stau. The degeneracy of the selectron and smuon guarantees that all constraints from flavor-changing neutral processes are satisfied.

In this simple model, there is also a shift in the mass of one of the left-handed squark doublets,

$$\frac{\delta m^2_{Q_3}}{m^2_Q} = -\frac{1}{16}\frac{y^2_d}{\alpha^2_3}\frac{|F_S|^2}{M^4} \approx -6 y^2_d \frac{|F_S|^2}{M^4} \tag{3.5}$$

[where $y^2_d = \Sigma_i (Y^d_{i4})^2$], and a related shift in the mass of the down Higgs boson, so that Eq. (3.4) is modified to

$$\frac{\delta m^2_{H_D}}{m^2_{H_D}} = -\frac{1}{9}\frac{3y^2_d + y^2_{\ell}}{\alpha^2_2}\frac{|F_S|^2}{M^4} \approx -10^2 (3y^2_d + y^2_{\ell})\frac{|F_S|^2}{M^4}. \tag{3.6}$$

Again, if the exotic Yukawa coupling is largest for Q_3, then Q_1 and Q_2 remain degenerate and constraints from flavor-changing neutral processes (e.g., K-$\bar{K}$ and D-$\bar{D}$ mixing) are easily satisfied. If we adopt $|F_S|/M^2 \lesssim 0.03$, then both Eqs. (3.5) and (3.6) are small.

The most plausible effect of the mixing is then a (negative) shift in the mass of $\tilde{\tau}_R$. There is a small shift in the squared mass of H_D, while for all other scalars, the one-loop mixing contribution is either absent or very small. It is possible, however, that $y_{\ell} \ll y_d \lesssim 1$. In this case, a substantial shift in $m^2_{H_D}$ is possible with a corresponding (but much less substantial) shift in $m^2_{Q_3}$. Finally, if the generation hierarchy of the exotic Yukawa couplings is very different from the ordinary Yukawa couplings, the result could be that, say, the selectron or the smuon is the lightest among the right-handed sleptons. But then the constraints from $\mu \to e\gamma$ are significant and require that the splitting be small. Similarly, constraints from K-$\bar{K}$ mixing require that the splitting in the squark sector be small if the largest exotic Yukawa coupling is Y^d_{14} or Y^d_{24}.

Next, consider models with $N_5 > 1$. Here, for generic mixing between messenger and matter fields, all three generations of left-handed squarks and of right-handed sleptons are split. Flavor-changing neutral current constraints are significant. But if we take, as above, $|F_S|/M^2 \lesssim 0.03$ and, in addition, assume that the exotic Yukawa couplings are not larger than the corresponding ordinary Yukawa couplings, e.g., $Y'_{4\mu} \lesssim m_\mu \tan\beta/m_t$, then all the constraints are satisfied. Such a hierarchy in the exotic Yukawa couplings is very likely if the smallness and hierarchy of the ordinary Yukawa couplings is explained by horizontal symmetries (see, for example, [22]).

Finally, we may consider mixing with messenger $\mathbf{10}+\overline{\mathbf{10}}$. Then masses of all ordinary scalar fields, except for the right-handed sleptons, are shifted. Again, for a single pair of $\mathbf{10}+\overline{\mathbf{10}}$, only one generation in each sector is affected. Very plausibly, these are the third generation sfermions, so that constraints from flavor-changing neutral current processes are rather weak. A small parameter $|F_S|/M^2$ guarantees that these one loop corrections are smaller than or comparable to the two-loop gauge contributions. Substantial corrections could occur for the slepton and Higgs fields, but the mass shifts for squarks are small.

We learn then that there are a few possibilities concerning the effects of messenger-matter mixing.

(a) There is no mixing or the mixing is negligibly small. Equations (1.2) and (1.3) remain valid. This is the situation if there is a symmetry that forbids mixing or if the ratio $|F_S|/M^2$ is small.

(b) There is a large negative mass shift of order one for $\tilde{\tau}_R$ and a small negative mass shift of order 0.1 for H_D. For all other soft supersymmetry-breaking parameters, Eqs. (1.2) and (1.3) remain an excellent approximation. This is the situation if $y_{\ell'}^2|F_S/M^2| \sim 0.03$.

(c) There is a large negative mass shift of order one for H_D and a small negative mass shift of order 0.02 for Q_3. For all other soft supersymmetry-breaking parameters, Eqs. (1.2) and (1.3) remain an excellent approximation. This is the situation if $y_{d'}^2|F_S/M^2| \sim 0.06$ and $y_{\ell'} \ll y_d$.

(d) The lightest squark or slepton could belong to the first or second generation or all three generations could be split in masses. This is the situation if the hierarchy in the exotic Yukawa couplings is different from that of the ordinary ones or if there are several $\mathbf{5}+\overline{\mathbf{5}}$ representations. But then phenomenological constraints require that the mass shifts be small.

We emphasize that the effects cannot be large in the squark sector. But there could be large effects in the slepton and/or Higgs sectors. Such corrections might be helpful in understanding at least one issue. In low energy breaking, there are potential fine-tuning problems in obtaining a suitable breaking of SU(2)×U(1). The problem is that the masses of the lightest right-handed leptons are constrained, from experiment, to be greater than about 45 GeV. On the other hand, if gauge mediation is the principle source of all masses, the contribution to the masses of the Higgs doublets tends to be larger. So if the lightest slepton has a mass of order 80 GeV or more (as suggested by the CDF event), then the typical contributions to Higgs boson masses would seem to be on the large side. Additional negative contributions would tend to ameliorate this problem.

Finally, we should mention another possible source of the violation of universality. Throughout this discussion, we have assumed an underlying R parity, and that q and $\bar{q}$ have the same R parity as ordinary quarks. It is possible that R parity is broken or that q and $\bar{q}$ have the *opposite* R parity. In this case, operators like $LQ\bar{q}$ or $\bar{u}\bar{d}\bar{q}$ may be allowed. The latter can lead to more appreciable shifts in squark masses and, thus, more significant violations of universality in the squark sector than we have contemplated up to now.

IV. μ PROBLEM

In the MGM model of Ref. [2], the following mechanism to generate a μ term was employed. An additional singlet field T was introduced, which couples to the Higgs fields through a nonrenormalizable term in the superpotential:

$$\frac{T^n}{M^{n-1}} H_U H_D. \tag{4.1}$$

To generate a B term, it was suggested that a term in the superpotential of the form

$$\lambda_h S H_U H_D \tag{4.2}$$

is allowed. With a small $\lambda_h \sim (\alpha_2/\pi)^2$, it gives an acceptable $B \sim (\alpha_2/\pi)^2 F_S$ and a negligible contribution to μ. It is difficult, if not impossible, to find a symmetry that forbids all $H_U H_D$ couplings except Eq. (4.1) and, with an appropriately small λ_h, Eq. (4.2). (For previous, unsuccessful attempts, see [14].) If, however, Eq. (4.2) is forbidden or highly suppressed, so that $B=0$ at the tree level, then loop contributions still generate $B \sim (\alpha_2/\pi)^2 \Lambda\mu$ [9], which is small, but not negligibly small. We now present a simple model where, indeed, as a result of a discrete symmetry, Eq. (4.1) gives the largest contribution to μ while λ_h of Eq. (4.2) is negligibly small. As we will explain in Sec. VI, such a model offers hope of solving the supersymmetric CP problem.

Let us introduce a (horizontal) symmetry $H=Z_m$ and set the H charges of the relevant fields to

$$H(S)=0, \quad H(H_U H_D)=n, \quad H(T)=-1. \tag{4.3}$$

The various VEV's are hierarchical, $\langle T\rangle \gg \sqrt{\langle F_S\rangle} \gg \langle H_U\rangle$, $\langle H_D\rangle$, and spontaneously break, respectively, the symmetry H, supersymmetry (and an R symmetry), and the electroweak symmetry. The relevant terms in the superpotential are

$$W = W_0(S) + W_1(S,T) + W_2(S,T,H_U,H_D),$$

$$W_1 \sim \frac{T^m}{M_P^{m-3}} \left(1 + \frac{S}{M_P} + \cdots\right),$$

$$W_2 \sim \frac{T^n H_U H_D}{M_P^{n-1}} \left(1 + \frac{S}{M_P} + \cdots\right). \tag{4.4}$$

Here, M_P is the Planck scale which suppresses all nonrenormalizable terms. The dots stand for terms that are higher order in S/M_P.

Similarly to the model of Abelian horizontal symmetries presented in [22], the minimum equations give an H-breaking scale that is intermediate between the supersymmetry breaking scale and the Planck scale and depends only on m. Explicitly, $\partial V/\partial T = 0$ gives

$$\frac{F_S}{M_P^2}\sim\left(\frac{T}{M_P}\right)^{m-2}. \tag{4.5}$$

Also similarly to the models of [22], the supersymmetric μ problem is solved because a term $\mu H_U H_D$ violates H. The leading contribution to μ is of order

$$\frac{\mu}{M_P}\sim\left(\frac{T}{M_P}\right)^{n}\sim\left(\frac{F_S}{M_P^2}\right)^{n/(m-2)}. \tag{4.6}$$

For definiteness, we take $F_S/M_P^2\sim 10^{-28}$ and require that Eq. (4.6) predict $\mu/M_P\sim 10^{-16}$. This is the case for $n\approx\frac{4}{7}(m-2)$. The simplest option is then $n=4$ and $m=9$ (corresponding to $T/M_P\sim 10^{-4}$). If one insists on larger T/M_P, so that it may be relevant to the fermion mass hierarchy, say 10^{-3} (10^{-2}), it can be achieved with $n=5$, $m=11$ ($n=8$, $m=16$).

A B term is also generated by W of Eq. (4.4). The leading contribution is of order

$$B\sim\frac{F_S\mu}{M_P}. \tag{4.7}$$

This contribution to B is $\ll\mu^2$ and, therefore, negligible. A much larger contribution is generated at the two-loop level (note that M_2 is generated by one loop diagrams):

$$B\sim\frac{\alpha_2}{\pi}\,\mu M_2. \tag{4.8}$$

This is smaller that the square of the electroweak symmetry-breaking scale by a factor of order α_2. Consequently, tan β is large, of order α_2^{-1}.

V. NATURALLY LARGE tan β

It has been argued [16–18] that, if there are only two Higgs doublets in the low energy supersymmetric model, large tan β requires a fine-tuning in the parameters of the Lagrangian of order (1/tanβ). The naturalness criterion used, for example, in Ref. [16] states that ''unless constrained by additional approximate symmetries, all mass parameters are about the same size, and all dimensionless numbers are of order one.'' However, in all existing models of dynamical supersymmetry breaking (DSB), there is more than one relevant energy scale. The assumption that all dimensionful parameters are characterized by a single scale may fail. Then large tanβ may arise naturally, as is the case in the model of the previous section.

Let us first repeat the argument that large tanβ requires fine-tuning. The basic assumption here is that, in the low energy effective supersymmetric standard model, there is a single scale that is the electroweak (or, equivalently, the supersymmetry) breaking scale. A dimensionful parameter can be much smaller only as a result of an approximate symmetry. The Higgs potential for the two Higgs doublets is

$$m_U^2H_U^2+m_D^2H_D^2+B(H_UH_D+\text{H.c.})+\frac{g^2+g'^2}{8}\,(|H_U|^2-|H_D|^2)^2. \tag{5.1}$$

In the large tanβ region,

$$\frac{1}{\tan\beta}\approx-\frac{B}{m_U^2+m_D^2}. \tag{5.2}$$

Large tanβ requires $B\ll m_U^2+m_D^2$. There are two symmetries that could suppress B below its natural value of order m_Z^2. If B is made small ($B\sim m_Z^2/\tan\beta$) by an approximate R symmetry, the W-ino mass M_2 should also be small ($M_2\sim m_Z/\tan\beta$). If B is made small by an approximate Peccei-Quinn (PQ) symmetry, then the μ term should also be small ($\mu\sim m_Z/\tan\beta$). This has interesting consequences for the chargino mass matrix:

$$\begin{pmatrix} \mu & \frac{g}{\sqrt{2}}\langle H_U\rangle \\ \frac{g}{\sqrt{2}}\langle H_D\rangle & M_2\end{pmatrix}. \tag{5.3}$$

As $\langle H_D\rangle$ is small by assumption and as (to make B naturally small) at least one of μ and M_2 has to be small, the mass matrix (5.3) leads to a light chargino (with mass of order $m_Z/\tan\beta$). This is phenomenologically unacceptable (the bounds on chargino masses are roughly $\gtrsim m_Z/2$). This means that the natural scale for either μ or M_2 is of the order of m_Z tanβ, and the criterion for naturalness is violated.

The assumption that a natural effective low energy supersymmetric standard model has a single energy scale is a strong one. In all existing models of DSB, there are at least three energy scales: the Planck scale M_P, the supersymmetry-breaking scale M_S and the electroweak breaking scale m_Z. Whether indeed m_Z is the only relevant scale for the low energy theory and, in particular, for μ and B, is a model-dependent question. In hidden sector models of supersymmetry breaking, one assumes that $m_Z\sim M_S^2/M_P$ is, indeed, the only relevant scale in the low energy model. But this is a rather arbitrary (though convenient) ansatz and, in the absence of a detailed high energy theory for the messenger sector, does not stand on particularly firm grounds. The situation is even more complicated in models of gauge-mediated supersymmetry breaking. Here, in addition to the Planck scale, there exist the dynamical supersymmetry-breaking scale $\sqrt{F_S}$, the scale $\Lambda=F_S/S$, and the electroweak scale $m_Z\sim\alpha_2\Lambda$. Which of these scales is relevant to B depends on the mechanism that generates B. It could very well be that the natural scale for B is $B\ll m_Z$.

To understand the situation in more detail, let us assume that there is neither a PQ symmetry nor an R symmetry to suppress B. Then the natural value for μ is M_P, and the model does not provide any understanding of the μ problem. But even if we assume that $\mu\ll M_P$ for some reason, the term (4.2) is allowed. This leads to $\mu\sim S$ and $B\sim F_S$. Both values are unacceptably large, but our main point here is that the natural scale for B could easily be the *highest* supersymmetry-breaking scale in the *full* theory, M_S.

In the model presented in the previous section, the H symmetry leads to an accidental PQ symmetry. The small breaking parameter of the PQ symmetry is of order m_Z/M_P, thus solving the μ problem. At the same time, it leads to a tree level value for B that is of order $F_S\mu/M_P$. This is actually similar to the scale in supergravity models, except that in those models $F_S\mu/M_P\sim m_Z^2$, while in models of GMSB $F_S\mu/M_P\ll m_Z^2$. Consequently, this contribution to B is negligibly small. The main point here is that the natural scale for B could be M_S^2/M_P; this scale coincides with the electroweak scale m_Z only in supergravity models.

Finally, a larger contribution to B arises in our model from two loop diagrams, of order $\alpha_2\mu M_2 \sim \alpha_2^2\mu\Lambda$. This is smaller than the electroweak scale $m_Z^2 \sim \alpha_2^2\Lambda^2$ by a factor of the order of $\mu/\Lambda \sim \alpha_2$. We learn that different combinations of scales could be relevant to B and to m_Z. If the combinations are such that $B \ll m_Z^2$, then a large $\tan\beta$ arises and no fine-tuning is required. The model of the previous section provides a specific example of this situation.

VI. SUPERSYMMETRIC *CP* PROBLEM

Supersymmetric theories introduce new sources of CP violation. With the minimal supersymmetric extension of the standard model and assuming universality of gaugino and of sfermion masses, there are four additional phases beyond the Kobayashi-Maskawa phase and $\theta_{\rm QCD}$ of the standard model. One phase appears in the μ parameter, and the other three in the soft supersymmetry-breaking parameters M_λ, A, and B:

$$\mathcal{L} = \tfrac{1}{2} M_\lambda \lambda\lambda - A(h_u Q H_U \overline{u} - h_d Q H_D \overline{d} - h_\ell L H_d \overline{e}) - B H_U H_D + \text{H.c.}, \tag{6.1}$$

where λ are the gauginos and h_i the Yukawa couplings. Only two combinations of the four phases are physical [23,24]. These can be taken to be

$$\phi_A = \arg(A^* M_\lambda),$$

$$\phi_B = \arg(B_\mu^* M_\lambda^*). \tag{6.2}$$

Unless these phases are $\lesssim 10^{-2}$ or supersymmetric masses are $\gtrsim 1$ TeV, the supersymmetric contribution to the electric dipole moment of the neutron is well above the experimental bound. This is the supersymmetric CP problem.

In models of GMSB, gaugino masses are not universal [see Eq. (1.2)]. However, with a minimal messenger sector ($N_S = N_5 = 1$), gaugino masses carry a universal phase. Thus there still exist only the two new phases defined in Eq. (6.2).

In the MGM model of Ref. [2], $A(\Lambda) = 0$. In its minimal version investigated in Ref. [9], also $B(\Lambda) = 0$. Radiative corrections give [9]

$$A_t \simeq A_q(\Lambda) + M_2(\Lambda)[-1.85 + 0.34|h_t|^2],$$

$$\frac{B}{\mu} \simeq \frac{B}{\mu}(\Lambda) - \frac{1}{2}A_t(\Lambda) + M_2(\Lambda)[-0.12 + 0.17|h_t|^2]. \tag{6.3}$$

Using Eq. (6.2), we learn from Eq. (6.3) that, for $A(\Lambda) = B(\Lambda) = 0$, one has

$$\phi_A = \phi_B = 0. \tag{6.4}$$

Thus the supersymmetric CP problem is solved in this model.

The vanishing of the supersymmetric phases goes beyond the approximation (6.3). It is actually common to all models with universal sfermion masses and a universal phase in the gaugino masses and where, at the tree level,[8] $A = B = 0$. In the absence of nongauge interactions, there is an additional R symmetry in the supersymmetric standard model. In a spurion analysis, it is possible to assign the same R charge to M_λ, A, and B [25,24]. If the only source of R symmetry breaking is gaugino masses, both ϕ_A and ϕ_B are zero, just because A, B, and the gaugino mass have the same R charge, and the RG evolution formally respects the R symmetry.

At the two-loop level, Yukawa interactions affect the running of A. Proportionality of the A terms and the Yukawa terms is violated and complex phases (related to the Kobayashi-Maskawa phase) appear in off-diagonal A terms (see Refs. [26, 27] for the relevant RGE). The contribution of these phases to the electric dipole moment of the neutron is, however, highly suppressed.

We conclude then that in the minimal version of MGM models (namely, when $A = B = 0$ at a high scale) the supersymmetric CP problem is solved.

VII. CONCLUSIONS

If more events with two photons plus missing energy are discovered, this can be viewed as strong evidence for low energy supersymmetry breaking. The MGM has a strong appeal, given its simplicity, but one can easily imagine that the messenger sector may be more complicated. It is possible that the data will support the MGM, but even given the limited information we have now, there are hints that some extension of the model may be required [7]. We have seen that in weakly coupled theories the spectrum can be modified in two significant ways. First, the hierarchy may be altered. As a result, one can imagine that, say, slepton doublets are not much more massive than singlets (as suggested in Ref. [7]). Second, there can be departures from universality. In other words, some SU(2) singlet sleptons might be lighter than others. We have seen that there are significant constraints on such universality violations coming, for example, from requiring reasonable breaking of SU(2)×U(1). We have also seen that if horizontal symmetries are responsible for the hierarchies of ordinary quark and lepton masses, at most only a few states will exhibit appreciable universality violation (e.g., the stau may be significantly lighter than the other sleptons).

We have so far avoided the more difficult question of what may happen in strongly coupled theories. These issues were touched upon in [4]. In the event that the underlying supersymmetry-breaking theory is strongly coupled, it seems likely that some of our constraints will be relaxed. For example, it is not clear that asymptotic freedom is a correct criterion, since we know from the work of Seiberg [28] that the infrared degrees of freedom of a theory may be quite different than the microscopic degrees of freedom. Another difficulty lies in mass formulas such as Eqs. (1.3) and (1.2). It is not clear whether in strongly coupled theories, the factors of $(4\pi)^{-2}$ which appear in weak coupling will also appear. For example, there may be single-particle states which can appear in a two-point function relevant to the gaugino mass computation, and one might suspect that the result,

[8]We thank Riccardo Rattazzi for explaining this point to us.

lacking the usual phase space factors, will be larger. Thus one can imagine that the SUSY-breaking scale might be closer than suggested by weak coupling models. This possibility should be taken seriously, since one might hope in such a framework to avoid the division into different sectors which we have seen is inevitable in weakly coupled models.

We have also discussed the μ problem and the question of large $\tan\beta$. We have noted that the usual arguments that large $\tan\beta$ requires fine-tuning make assumptions about the scales μ and B which need not hold—indeed one might argue are not likely to hold—in theories of low energy dynamical breaking. In particular, it is quite natural for B to be very small at the high scale [9]. In this situation, the supersymmetric CP problem is automatically solved.

The MGM models are attractive in that they are highly predictive, guarantee universality, can suppress the supersymmetric CP-violating phases, and predict events with final photons and missing energy similar to the one observed by the CDF. In this work we have learned that reasonable extensions of the minimal models retain many of these nice features while offering a richer phenomenology.

(a) The number of parameters describing sfermion and gaugino masses can increase to 8 with extended messenger sectors or to about 11 with messenger-matter mixing. The hierarchy of masses between, say, gauginos and sfermions or squarks and sleptons may be different from the minimal models.

(b) Universality is violated with messenger-matter mixing, but most likely, it is only the third generation that is significantly affected. Interesting flavor-changing neutral current processes may be observed, for example, in τ decays.

(c) Final photons and missing energy remain the typical signature of low energy supersymmetry breaking, but the detailed nature of the final states could be rather different than in the MGM models.

ACKNOWLEDGMENTS

We thank Savas Dimopoulos, Scott Thomas, Francesca Borzumati, and Riccardo Rattazzi for important comments and criticisms as these ideas were developed. Y.N. was supported in part by the United States–Israel Binational Science Foundation (BSF), by the Israel Commission for Basic Research, and by the Minerva Foundation (Munich). The work of M.D. was supported in part by the U.S. Department of Energy.

[1] M. Dine, A. E. Nelson, and Y. Shirman, Phys. Rev. D **51**, 1362 (1995).
[2] M. Dine, A. E. Nelson, Y. Nir, and Y. Shirman, Phys. Rev. D **53**, 2658 (1996).
[3] CDF Collaboration, S. Park, in *10th Topical Workshop on Proton-Antiproton Collider Physics*, edited by R. Raha and J. Yoh, AIP Conf. Proc. No. 357 (AIP, New York, 1995).
[4] S. Dimopoulos, M. Dine, S. Raby, and S. Thomas, Phys. Rev. Lett. **76**, 3494 (1996).
[5] S. Ambrosanio, G. L. Kane, G. D. Kribs, S. P. Martin, and S. Mrenna, Phys. Rev. Lett. **76**, 3498 (1996).
[6] D. R. Stump, M. Wiest, and C. P. Yuan, Phys. Rev. D **54**, 1936 (1996).
[7] S. Dimopoulos, S. Thomas, and J. D. Wells, Phys. Rev. D **54**, 3283 (1996).
[8] S. Ambrosanio, G. L. Kane, G. D. Kribs, S. P. Martin, and S. Mrenna, Phys. Rev. D **54**, 5395 (1996).
[9] K. S. Babu, C. Kolda, and F. Wilczek, Phys. Rev. Lett. **77**, 3070 (1996).
[10] L. E. Ibanez and D. Lust, Nucl. Phys. **B382**, 305 (1992).
[11] V. Kaplunovsky and J. Louis, Phys. Lett. B **306**, 269 (1993).
[12] T. Banks and M. Dine, Nucl. Phys. **B479**, 173 (1996).
[13] S. Dimopoulos, G. F. Giudice, and A. Pomarol, ''Dark Matter in Theories of Gauge-Mediated Supersymmetry Breaking,'' Report No. CERN-TH/96-171, hep-ph/9607225 (unpublished).
[14] T. Gherghetta, G. Jungman, and E. Poppitz, ''Low-Energy Supersymmetry Breaking and Fermion Mass Hierarchies,'' Report No. hep-ph/9511317 (unpublished).
[15] G. Dvali, G. F. Giudice, and A. Pomarol, ''The Mu Problem in Theories with Gauge Mediated Supersymmetry Breaking,'' Report No. hep-ph/9603228 (unpublished).
[16] A. E. Nelson and L. Randall, Phys. Lett. B **316**, 516 (1993).
[17] L. J. Hall, R. Rattazzi, and U. Sarid, Phys. Rev. D **50**, 7048 (1994).
[18] R. Rattazzi and U. Sarid, Phys. Rev. D **53**, 1553 (1996).
[19] S. Dimopoulos and H. Georgi, Nucl. Phys. **B193**, 150 (1981).
[20] A. E. Farraggi, Phys. Lett. B **387**, 775 (1996).
[21] T. Hotta, K.-I. Izawa, and T. Yanagida, ''Dynamical Supersymmetry Breaking without Messenger Gauge Interactions,'' Report No. hep-ph/9606203 (unpublished).
[22] M. Leurer, Y. Nir, and N. Seiberg, Nucl. Phys. **B420**, 468 (1994).
[23] M. Dugan, B. Grinstein, and L. Hall, Nucl. Phys. **B255**, 413 (1985).
[24] S. Dimopoulos and S. Thomas, Nucl. Phys. **B465**, 23 (1996).
[25] L. J. Hall, R. Rattazzi, and U. Sarid, Phys. Rev. D **50**, 7048 (1994).
[26] S. Bertolini, F. Borzumati, A. Masiero, and G. Ridolfi, Nucl. Phys. **B353**, 591 (1991).
[27] S. Bertolini and F. Vissani, Phys. Lett. B **324**, 164 (1994).
[28] N. Seiberg, Nucl. Phys. **B435**, 129 (1995).