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
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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 Λ and B vanish at the tree level, tanβ could be large without fine-tuning. We explain how the supersymmetric CP problem is solved in such models. [S0556-2821(97)03603-5]

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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^− γ γ E_T event observed at the Collider Detector at Fermilab (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 5 + 5 of SU(5), i.e., color triplets, q + ¯q, and weak doublets, l + l. These couple to a single gauge-singlet field S through a superpotential

$$W = λ_1 S q ̅ q + λ_2 S l ̅ l.$$  (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

$$m_λ = c_i \frac{α_i}{4 \pi} Λ,$$  (1.2)

where Λ = F_S / S, c_1 = 5/3, c_2 = c_3 = 1, and α_i = α_i / cos^2 θ_W. For the scalar masses one has

$$m^2 = 2 Λ^2 \left[ C_3 \left( \frac{α_3}{4 π} \right)^2 + C_2 \left( \frac{α_2}{4 π} \right)^2 + \frac{5}{3} \left( \frac{Y}{2} \right)^2 \left( \frac{α_1}{4 π} \right)^2 \right],$$  (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.\(^4\)

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.\(^5\) On the other hand, while

\(^4\)These formulas predict a near degeneracy of the B-parity and the right-handed sleptons. Important corrections due to operator renormalization and D terms have been discussed in [9].

\(^5\)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. 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 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.

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In this paper, we will use the term ‘‘universality’’ to mean scalar masses which are functions only of gauge quantum number and $\Lambda$ terms which are small or proportional to fermion Yukawa couplings.

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.

naturalness in more detail. We present a specific version of the minimal model of Ref. [9] where a discrete symmetry predicts $\mu$ 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 $\phi_A$ and $\phi_B$ 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, 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

5This 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 renor-
malizable interactions) to the messengers.\textsuperscript{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 uni-
fication with a desert, they must come in complete SU(5)
multiplets. (For a different scenario, see Ref. [20].) More-
over, 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).\textsuperscript{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 modifica-
tions, 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 in-
volving superpotential couplings or gauge interactions. Aes-
thetic issues aside, such models will have difficulty explain-
ing 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
$\mu$ 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 $N_5$
of singlet fields with nonvanishing scalar and $F$ components.
If there are $N_5$ $\mathbf{5} + \overline{\mathbf{5}}$'s and $N_5$ 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

\[ \sum_{i=1}^{N_5} \sum_{j,K=1}^{N_5} (\lambda_1)_{iJK} S_i \bar{q} \bar{f} + (\lambda_2)_{iJK} S_i T K. \] (2.1)

For large enough $N_5$ and $N_6$, the mass 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
$N_5=2$ and $N_5=1$. Suppose that $S_1$ has a large scalar compo-
nent 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 sleptons comparable
supersymmetry-conserving masses, while couplings to $S_2$
could be, say, order 1 for messenger quarks, but small for
messenger sleptons. This could alter the hierarchy (in a uni-
versal 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 hier-
archy is altered.

On the other hand, the modifications of the hierarchy can-
not 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 ex-
tremely 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 $\gamma \gamma$ events in this framework, the fundamental scale of super-
symmetry breaking cannot be much larger than $10^{13}$ to
$10^{14}$ 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 hier-
archy might arise, but not violations of universality. The
fields $\bar{q}$ have the same quantum numbers as the ordinary
$d$ fields. We define $\bar{d}$ 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, $H_d Q \bar{d}$. Indeed,
terms of the form $H_d Q \bar{g}$ can be forbidden by discrete sym-
metries. (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 cosmol-
yogy. In the next section, we explore the consequences of
introducing such couplings.

\section{III. MESSENGER-MATTER MIXING}

In order to understand possible modifications of the spec-
trum 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 sleptons, while $\langle F \rangle$
leads to a supersymmetry-violating splitting in these multiplets. At

\footnotesize
\textsuperscript{6}We are assuming here that $R$ parity is conserved, but all of the

\textsuperscript{7}Alternatively, there might be “bare mass” in the superpoten-
tial analogous to the $\mu$ term. These might arise by the mechanism
for the $\mu$ term described in Ref. [2] and, further, in Sec. IV.
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 $\Lambda = F_S/S$ (here below, expectation values are understood).

The simple modification that we consider takes place in the Yukawa sector. The messenger $/\bar i$ field has the same gauge quantum numbers as the ordinary lepton doublets; at one loop, gauginos gain mass through their couplings to these fields; at two loops, ordinary squarks and sleptons gain mass. Thus, in the absence of a symmetry, one expects these fields to mix. In particular, in the Yukawa couplings

$$H_D L_i Y_i e_{\bar{i}} + H_D Q_i Y_i d_{\bar{i}}, \quad (3.1)$$

each of $L_i$ and $d_{\bar{i}}$, refers to the four objects with the same quantum numbers. Then $Y'$ is a $4\times3$ matrix while $Y^d$ is a $3\times4$ matrix. By convention, we call $L_4$ and $d_{\bar{4}}$ the linear combination of fields which couple to $S$ in Eq. (1.1). We refer to $3'$ and $3_{\bar{4}}$ 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 $F_S/S$. 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, $\tilde e_3$, has a substantial Yukawa coupling to $L_4$ and call this coupling $y'_{3}$. [There is actually no loss of generality here. In general, the affected slepton is the combination $\Sigma_i Y'_{4i}\tilde e_i$, with $y'_{3} = \Sigma_i (Y'_{4i})^2$.] The mass shift is

$$\frac{\delta m^2_{\tilde e_3}}{m_{\tilde e_3}} = - \frac{M^2}{6} \frac{y'_{3}^2 |F_S|^4}{16 \alpha^2} \frac{|F_S|^2}{M^4}. \quad (3.2)$$

where $M = \lambda S$ is the mean mass of the $\tilde e_3$ multiplet. Using Eq. (1.3) for the (universal) two-loop contribution to $m_{\tilde e_3}$, we find

$$\frac{\delta m^2_{\tilde e_3}}{m_{\tilde e_3}} = - \frac{1}{12} \frac{y'_{3}^2 |F_S|^2}{\alpha_y} \frac{|F_S|^2}{M^4} \approx - 10^3 y_{3}^2 \frac{|F_S|^2}{M^4}. \quad (3.3)$$

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

$$\frac{\delta m^2_{H_D}}{m_{H_D}} = - \frac{1}{9} \frac{y_{3}^2 |F_S|^2}{\alpha_y} \frac{|F_S|^2}{M^4} \approx - 10^3 y_{3}^2 \frac{|F_S|^2}{M^4}. \quad (3.4)$$

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, $\delta m^2_{\tilde e_3}/m_{\tilde e_3} \approx 1$ if $|F_S|^2/M^2 < 0.03$.

(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 $5 + \bar 5$ pair, this mass shift affects only one right-handed slepton generation. The other two remain degenerate.

(iv) If, similarly to ordinary Yukawa couplings, $Y'_{\bar 4i} \gg Y'_{4i}$ 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,

$$\frac{\delta m^2_{Q_3}}{m_{Q_3}} = - \frac{1}{16} \frac{y_{3}^2 |F_S|^2}{\alpha_y} \frac{|F_S|^2}{M^4} \approx - 6 y_{3}^2 \frac{|F_S|^2}{M^4}. \quad (3.5)$$

[where $y_{3}^2 = \Sigma_i (Y_{3i})^2$, and a related shift in the mass of the down Higgs boson, so that Eq. (3.4) is modified to]

$$\frac{\delta m^2_{H_D}}{m_{H_D}} = - \frac{1}{9} \frac{3 y_{3}^2 + y_4^2 |F_S|^2}{\alpha_y} \frac{|F_S|^2}{M^4} \approx - 10^2 (3 y_{3}^2 + y_4^2) \frac{|F_S|^2}{M^4}. \quad (3.6)$$

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., $K-K$ and $D-D$ mixing) are easily satisfied. If we adopt $|F_S|^2/M^2 < 0.03$, 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 $\tilde e_3$. 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 $y_{3}^2 \ll y_2^2 \approx 1$. In this case, a substantial shift in $m_{H_D}$ is possible with a corresponding (but much less substantial) shift in $m_{Q_3}$. 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 $K-K$ mixing require that the splitting in the squark sector be small if the largest exotic Yukawa coupling is $Y_{4\bar 4}$ or $Y_{4\bar 4}^3$.

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 signifi-
cant. But if we take, as above, $|F_3|/M^2 \leq 0.03$ and, in addition, assume that the exotic Yukawa couplings are not larger than the corresponding ordinary Yukawa couplings, e.g., $y'/\tan \beta \approx m_t$, 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 $10 + 10^c$. Then masses of all ordinary scalar fields, except for the right-handed sleptons, are shifted. Again, for a single pair of $10 + 10^c$, 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 $|F_3|/M^2$ 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 $|F_3|/M^2$ is small.

(b) There is a large negative mass shift of order one for $\tilde t_R$ 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 $y_u^2|F_3/M^2| \sim 0.03$.

(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 $y_d^2|F_3/M^2| \sim 0.06$ and $y_u \ll y_d$.

(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 $\bar 5 + 5$ 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 $SU(2) \times U(1)$. 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 $\tilde q$ have the same $R$ parity as ordinary quarks. It is possible that $R$ parity is broken or that $q$ and $\tilde q$ have the opposite $R$ parity. In this case, operators like $LQ\tilde q$ or $\bar u \bar d q$ 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. $\mu$ Problem**

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

$$\frac{T^n}{M^{n-1}} H_U H_D.$$  \hspace*{1cm} (4.1)

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

$$\lambda_b S H_U H_D$$  \hspace*{1cm} (4.2)

is allowed. With a small $\lambda_b \sim (\alpha_Z/\pi)^2$, it gives an acceptable $B \sim (\alpha_Z/\pi)^2 F_S$ and a negligible contribution to $\mu$. 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 $\lambda_b$, 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 $B \sim (\alpha_Z/\pi)^2 \mu$ [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 $\mu$ while $\lambda_b$ 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 $H = Z_n$ and set the $H$ charges of the relevant fields to

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

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

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

$$W_1 \sim \frac{T^n}{M^n_p} \left(1 + \frac{S}{M_p} + \cdots\right).$$

$$W_2 \sim \frac{T^n H_U H_D}{M^n_p} \left(1 + \frac{S}{M_p} + \cdots\right).$$  \hspace*{1cm} (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
Also similarly to the models of [22], the supersymmetric \( \mu \) problem is solved because a term \( \mu H_U H_D \) violates \( H \). The leading contribution to \( \mu \) is of order

\[
\frac{M_F}{M_p} \sim \left( \frac{T}{M_p} \right)^{m-2}.
\]

(4.5)

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

\[
\left( \begin{array}{cc}
\frac{\mu}{\sqrt{2}} \langle H_D \rangle & \frac{g}{\sqrt{2}} \langle H_U \rangle \\
\frac{g}{\sqrt{2}} \langle H_D \rangle & M_2
\end{array} \right).
\]

(5.3)

As \( \langle H_D \rangle \) is small by assumption and as (to make \( B \) naturally small) at least one of \( \mu \) and \( M_\tau \) has to be small, the mass matrix (5.3) leads to a light chargino (with mass of order \( m_\tau \tan \beta \)). This is phenomenologically unacceptable (the bounds on chargino masses are roughly \( \gg m_\tau / 2 \)). This means that the natural scale for either \( \mu \) or \( M_\tau \) is of the order of \( m_\tau \tan \beta \), 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 \( \mu \) and \( B \), is a model-dependent question. In hidden sector models of supersymmetry breaking, one assumes that \( m_Z \sim M_\tau^2 / M_p \) 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 \( \sqrt{F_\tau} \), the scale \( \Lambda = F_\tau / S \), and the electroweak scale \( m_Z / \alpha \Lambda \). 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 \( B \approx m_Z \).

To understand the situation in more detail, let us assume that there is neither a PQ symmetry nor an \( \tilde{R} \) symmetry to suppress \( B \). Then the natural value for \( \mu \) is \( M_p \), and the model does not provide any understanding of the \( \mu \) problem. But even if we assume that \( \mu \sim M_p \) for some reason, the term (4.2) is allowed. This leads to \( \mu \sim S \) and \( B \sim F_\tau \). 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 \( \mu \) problem. At the same time, it leads to a tree level value for \( B \) that is of order \( F_\tau \mu / M_p \). This is actually similar to the scale in supergravity models, except that in those models \( F_\tau \mu / M_p \sim m_Z \), while in models of GMSB \( F_\tau \mu / M_p \ll m_Z \). Consequently, this contribution to \( B \) is neg-

V. NATURALLY LARGE \( \tan \beta \)

It has been argued [16–18] that, if there are only two Higgs doublets in the low energy supersymmetric model, large \( \tan \beta \) requires a fine-tuning in the parameters of the Lagrangian of order (1/\( \tan \beta \)). 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 \( \tan \beta \) may arise naturally, as is the case in the model of the previous section.

Let us first repeat the argument that large \( \tan \beta \) 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

\[
m^2_U H_U^2 + m^2_D H_D^2 + B (H_U H_D + H.c.)
\]

\[
+ \frac{g^2 + g'^2}{8} (|H_U|^2 - |H_D|^2)^2.
\]

(5.1)

In the large \( \tan \beta \) region,
ligibly small. The main point here is that the natural scale for
B could be $M_{\Sigma}^2/M_{P}$; 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 $\alpha_2^2 \mu M^2_\Sigma \sim \alpha_2^2 \mu \Lambda$. 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 $\theta_{QCD}$ of the standard model. One phase appears in
the $\mu$ parameter, and the other three in the soft supersymmetry-breaking parameters
$M_\lambda$, $A$, and $B$:

$$\mathcal{L} = \frac{1}{2} M_\lambda \lambda - A(h_u Q H_u \tilde{u} - h_d Q H_d \tilde{d} - h_l L H \tilde{e})$$
$$- B H_u H_d + \text{H.c.}, \quad (6.1)$$

where $\lambda$ 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

$$\phi_A = \arg(A^* M_\lambda),$$
$$\phi_B = \arg(B^* M_\lambda). \quad (6.2)$$

Unless these phases are $\leq 10^{-2}$ or supersymmetric masses are $\geq 1$ TeV, the supersymmetric contribution to the
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
($N_3 = N_4 = 1$), 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]

$$A_i = A_i(\Lambda) + M_2(\Lambda)[-1.85 + 0.34|h_i|^2],$$
$$B \approx \frac{B}{\mu} \left( \frac{\Lambda}{\mu} \right) - \frac{1}{2} A_i(\Lambda) + M_2(\Lambda)[-0.12 + 0.17|h_i|^2]. \quad (6.3)$$

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

$$\phi_A = \phi_B = 0. \quad (6.4)$$

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,\footnote{We thank Riccardo Rattazzi for explaining this point to us.} $A = B = 0$. In
the absence of nongauge interactions, there is an additional $R$
symmetry in the supersymmetric standard model. In a spu-
tation analysis, it is possible to assign the same $R$ charge to
$M_\lambda$, $A$, and $B$ [25,24]. If the only source of $R$ symmetry
breaking is gaugino masses, both $\phi_A$ and $\phi_B$ are zero, just
because $A, B$, and the gaugino mass have the same $R$ charge,
and the $R G$ evolution formally respects the $R$ symmetry.

At the two-loop level, Yukawa interactions affect the run-
ing 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 super-
symmetric 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 ap-
peal, 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 lim-
ited information we have now, there are hints that some ex-
tension 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 re-
quiring reasonable breaking of SU(2)$\times$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 viola-
tion (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 ex-
ample, 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 fac-
tors of $(4 \pi)^{-2}$ which appear in weak coupling will also ap-
pear. 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,
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 $\mu$ 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 $\mu$ 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 $\tau$ 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.

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