



Some Remarks On Gauge-Mediated Supersymmetry Breaking

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Abstract

We investigate the communication of supersymmetry breaking to the Standard Model in theories of gauge mediated supersymmetry breaking with general weakly coupled messenger sectors. We calculate the one loop gaugino and two loop soft scalar masses for nonvanishing $Str M^2$ of the messenger sector. The soft scalar masses are sensitive to physics at scales higher than the messenger scale, in contrast to models with vanishing messenger supertrace. We discuss the implications of this ultraviolet sensitivity in theories with renormalizable and nonrenormalizable supersymmetry breaking sectors. We note that the standard relation, in minimal gauge mediation, between soft scalar and gaugino masses is altered in models with nonvanishing messenger supertrace.

1 Introduction.

In the past couple of years, models where supersymmetry breaking is communicated to the Standard Model by gauge interactions have received increasing attention [1]-[3]. This development has been stimulated in part by the new advances in understanding the dynamics of $N = 1$ supersymmetric gauge theories [4] and, recently, by the observation of the $e^+e^-\gamma\gamma$ event at Fermilab [5]. Gauge mediated supersymmetry breaking offers a predictive and testable alternative to the supergravity models and naturally suppresses flavor changing neutral currents.

To communicate supersymmetry breaking via gauge interactions one postulates the existence of heavy vectorlike multiplets of the Standard Model gauge group. These heavy “messenger” fields acquire soft supersymmetry breaking mass splittings due to their interactions with the supersymmetry breaking sector. Supersymmetry breaking is then transmitted to the Standard Model gauginos, which acquire mass at one loop level, and the squarks, sleptons and higgses, which acquire mass (squares) at the two loop level.

The interaction of the messenger fields (denoted by Q_1 and Q_2 , transforming in conjugate representations of the Standard Model gauge groups) with the supersymmetry breaking sector can be written in terms of a “spurion” field S —which is a dynamical field of the supersymmetry breaking sector—that acquires a supersymmetry breaking vacuum expectation value:

$$\langle S \rangle = s + \theta^2 F_s . \quad (1.1)$$

The most general interaction¹ between the spurion S and the messengers Q_i , quadratic in the messenger fields, can be written as [6]:

$$\int d^4\theta \frac{S^\dagger S}{M^2} f_i Q_i^\dagger \cdot Q_i + \left(\int d^2\theta S Q_1 \cdot Q_2 + \text{h.c.} \right) . \quad (1.2)$$

The scale M is a scale characteristic of the supersymmetry breaking sector. Inserting the expectation value (1.1) of the spurion generates soft scalar masses for the messengers and gives rise to the following general scalar mass matrix:

$$(Q_1^\dagger \ Q_2) \begin{pmatrix} a^2 & c^2 \\ c^{2*} & b^2 \end{pmatrix} \begin{pmatrix} Q_1 \\ Q_2^\dagger \end{pmatrix} , \quad (1.3)$$

while the fermion Dirac mass is equal to s . The scalar mass matrix in eq. (1.3) is the most general one can write for a single messenger multiplet. The elements a, b of the mass matrix are real, while c is complex; these four real parameters can be related to the parameters in eqs. (1.1, 1.2).

¹For simplicity, in eq. (1.2) we impose a (discrete) R symmetry to forbid terms like $S^\dagger Q_1 \cdot Q_2$, etc., in the Kähler potential. This does not affect the generality of the resulting mass matrix.

We note that if the coefficients f_i in eq. (1.2) are not small and if the scale M is not much larger than the vacuum expectation value $\langle S \rangle$ of the spurion, the nonholomorphic scalar masses—the elements a^2 and b^2 in (1.3)—receive soft supersymmetry breaking contributions from the D-terms, in addition to the (supersymmetric) contribution from the F-term. The supertrace of the messenger mass matrix, $\text{Str} M_{\text{mess}}^2 \equiv 2a^2 + 2b^2 - 4s^2$, is then nonvanishing.

In the models of Dine, Nelson, and Shirman (hereafter referred to as the “minimal gauge mediation”, MGM models, see ref. [2]), the coefficients f_i in eq. (1.2) are generated by loop effects² (while $M \simeq s$), and are therefore suppressed—the supertrace of the messengers’ squared mass matrix vanishes to a good accuracy. The models of ref. [2], however, require a rather complicated structure in order to give both a supersymmetry preserving and a supersymmetry breaking expectation value of the singlet field S . Moreover, since the messengers are not part of the supersymmetry breaking sector, the minimum with the required supersymmetry breaking expectation value of S is only local. Additional complications are needed in order to avoid this problem [7].

It appears natural, therefore, to look for models where the messengers are an intrinsic part of the supersymmetry breaking dynamics [8]. Recently several models of this type have been constructed [9], [10], [12]. The supersymmetry breaking sectors of these models are based in part on the $SU(N) \times SU(N - M)$ models, with $M = 1, 2$ [13]. At low energies, the gauge dynamics in these models can be integrated out. The infrared dynamics of the supersymmetry breaking sector is then described by a *weakly coupled* nonlinear supersymmetric sigma model, which contains the messenger fields (i.e. the fields Q_i above), as well as several gauge singlet fields that are essential for supersymmetry breaking³. These Standard Model gauge singlets play the role of the spurion S in eq. (1.2). The interaction between the “spurions” and the messengers can be written as in (1.2), with the scale M being identified with the scale of the vacuum expectation value, s , of S . The coefficients f_i are not loop suppressed and the supertrace of the messenger fields’ mass squared matrix is nonvanishing.

For most of the present investigation the detailed dynamics of these models is not important; we will only use them as an “existence proof” of models with weakly coupled messenger sectors with nonzero supertrace. It is natural to expect that in any dynamical model where the messengers participate in the supersymmetry breaking, and the low energy dynamics can be described by a weakly coupled nonlinear sigma model, $\text{Str} M_{\text{mess}}^2 \neq 0$. This can be seen by considering the tree level supertrace mass squared sum rule [11] for a general nonlinear sigma

²For example, the same two-loop graphs that generate soft masses for the ordinary squarks and sleptons also generate corrections to the messengers’ soft masses—the D-terms in eq. (1.2)

³In these models the Standard Model soft parameters receive contributions at energies higher than those in the sigma model as well—we will have more to say about these contributions in Section 3.

model:

$$\text{Str } M^2 = - 2 R_{ij^*} K^{il^*} K^{mj^*} W_m W_{l^*}, \quad (1.4)$$

where we use the notations of ref. [11]: W_m is the gradient of the superpotential, K^{ij^*} is the inverse Kähler metric, and R_{ij^*} is the Ricci tensor of the Kähler manifold. In eq. (1.4) the trace is taken over all states in the sigma model (including the messenger singlets), however—and explicit examples confirm this—one does not expect a restriction of the supertrace on the space of states charged under a global symmetry to vanish for a Kähler manifold with nonzero curvature.

In the next section we consider in detail the effect of the nonvanishing supertrace on the communication of supersymmetry breaking to the Standard Model fields. We find that the two loop scalar soft masses are sensitive to physics at momenta higher than the messenger scale, in contrast with the earlier models. This sensitivity can lead to an enhancement or suppression of the scalar soft masses compared to the gaugino masses. In Section 3, we consider the implications of this ultraviolet sensitivity on the calculability of the soft scalar masses in renormalizable and nonrenormalizable models of supersymmetry breaking. We point out the importance of possible threshold (matching) contributions of heavy states in the supersymmetry breaking sector that carry Standard Model quantum numbers. Finally, we point out the possibility of obtaining superparticle spectra with squarks and leptons lighter than the gauginos in the “hybrid” models of supersymmetry breaking. Many of the results reported in this letter have also been obtained by N. Arkani-Hamed, J. March-Russell, and H. Murayama, and reported in [10]—we thank these authors for discussions.

2 The two loop soft scalar masses with $\text{Str } \mathbf{M}_{\text{mess}}^2 \neq \mathbf{0}$.

In this section we will turn to a detailed calculation of the induced soft masses. Our starting point is the messenger mass matrix of eq. (1.3) which when diagonalized is of the form:

$$M_S^2 = \begin{pmatrix} a^2 & c^2 \\ c^{2*} & b^2 \end{pmatrix} = U \cdot \begin{pmatrix} m_1^2 & 0 \\ 0 & m_2^2 \end{pmatrix} \cdot U^\dagger, \quad (2.1)$$

where the unitary matrix U is

$$U = \begin{pmatrix} x & \sqrt{1-x^2} e^{i\alpha} \\ -\sqrt{1-x^2} e^{-i\alpha} & x \end{pmatrix}, \quad (2.2)$$

with $|x| \leq 1$. The Dirac mass of the messenger fermion will be denoted by m_f .

We note once again that the mass matrix above—unlike the MGM case—allows for the supertrace of the messengers to be nonvanishing⁴. This feature will in fact play a crucial role

⁴In case the messengers belong to several different representations the corresponding parameter is the

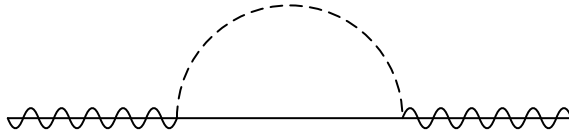


Figure 1: One loop messenger contribution to Standard Model gaugino masses.

in the discussion below. It is also worth mentioning that even for vanishing supertrace the above ansatz is more general than the one considered in refs. [17], [21].

The calculation of the radiatively induced gaugino and scalar masses is most easily performed in components. Superfield techniques [14] in theories with broken supersymmetry are generally useful for finding the infinite parts of Feynman diagrams—i.e. for calculating anomalous dimensions and beta functions—or for calculating finite parts when the mass splittings in the supermultiplets are small. It is in these cases only that the supersymmetry breaking effects can be treated as insertions in the relevant Feynman graphs, see e.g. ref. [15]. We are however interested in the more general case when the supersymmetry breaking splittings are not small compared to the supersymmetric messenger mass. Treating supersymmetry breaking effects simply as insertions is not appropriate in this case—one needs to use the exact superfield propagators in the supersymmetry breaking background. Expressions for the superfield propagators in a general supersymmetry breaking background have been given in the literature [16]. To the best of our knowledge, only the chiral superfield propagators in the supersymmetry breaking background have been obtained; even these are prohibitively complicated to allow for (two-) loop calculations.

We first consider the contributions to the Standard Model gaugino masses. These arise at one loop. The corresponding component graph, shown in Fig. 1, is finite and the calculation is a straightforward extension of that in [17]. The result for the Standard Model gaugino masses is:

$$m_a = e^{-i\alpha} x \sqrt{1-x^2} \frac{g_a^2 m_f}{4 \pi^2} S_Q \frac{y_1 \log y_1 - y_2 \log y_2 - y_1 y_2 \log(y_1/y_2)}{(y_1 - 1)(y_2 - 1)} \quad (2.3)$$

where $y_1 = m_1^2/m_f^2$, $y_2 = m_2^2/m_f^2$, g_a is the corresponding Standard Model gauge coupling, and S_Q is the Dynkin index of the messenger representation (normalized to 1/2 for a fundamental of $SU(N)$, while for $U(1)_Y$ it is simply Y^2 , where Y is the messenger hypercharge; we do not use GUT normalization of hypercharge). m_a above refers to the coefficient of the holomorphic gaugino bilinear operator $\lambda^{\alpha a} \lambda_{\alpha}^a$ in the Lagrangian.

We now turn to a consideration of the scalar masses. These arise at two loops⁵. As we

supertrace of the mass squared matrix weighted by the Dynkin index of the messenger representation. In the subsequent discussion we will continue to refer to this loosely as the supertrace of the mass matrix.

⁵ We note that with the more general mass matrix (2.1) one loop contributions to the hypercharge D-term are also possible. It is easy to see, however, that these are proportional to $2x^2 - 1$, and therefore negligible when

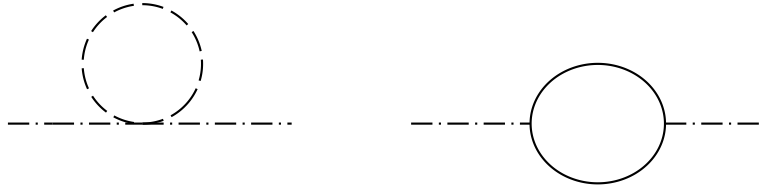


Figure 2: Divergent one-loop messenger contributions to the mass of the epsilon-scalars (dash-dotted lines), proportional to the supertrace of the messenger mass matrix.

will see below when the supertrace does not vanish the full contribution is in fact ultraviolet divergent. Uncovering this divergence though is subtle and requires a careful regularization of the theory.

We turn to this issue next. We will use dimensional reduction (DRED) [18] to regulate the theory in this paper. DRED insures that the supersymmetry Ward identities are preserved—at least to the two loop order of the calculation [19]. It also guarantees that the leading divergences cancel when all two loop graphs contributing to the scalar masses are summed up (if one uses conventional dimensional regularization instead, even for vanishing supertrace one finds that the divergences do not cancel—clearly a result of the fact that continuation to an arbitrary dimension does not preserve supersymmetry). Recall that in DRED one considers the theory *compactified* to $n = 4 - 2\epsilon$ dimensions, with $\epsilon > 0$. Now while the number of spinor components does not change under compactification, a vector field decomposes as an n -dimensional vector and 2ϵ scalar multiplets in the adjoint of the gauge group—the so-called epsilon-scalars. These 2ϵ scalar adjoint multiplets need to be fully incorporated in DRED for consistency (for a clear introduction and a discussion of the role of the epsilon-scalars see ref. [19]).

Note that in the theory with broken supersymmetry a mass term for the epsilon scalars is not forbidden by gauge invariance. In fact the epsilon scalars receive a divergent contribution to their mass at one loop from the graphs shown in Fig. 2. This divergent contribution is proportional to the supertrace of the messenger mass matrix [20] and is given by:

$$\delta m_\epsilon^2 = -S_Q \frac{g_a^2}{16\pi^2} \frac{\text{Str } M_{\text{mess}}^2}{\epsilon}, \quad (2.4)$$

where, as in (2.3), S_Q is the Dynkin index of the messenger representation. Correspondingly a one-loop counterterm needs to be added to fully renormalize the theory—we will choose the $x \simeq 1/\sqrt{2}$. Alternatively, if the messengers fall in complete $SU(5)$ multiplets, these contributions cancel for any value of x . For a generic $x \neq 1/\sqrt{2}$, however, one has to also worry about the two-loop contributions to the hypercharge D-term. Finally, the latter contributions can be controlled by imposing a discrete symmetry—one could have two sets of $SU(5)$ $\bar{5} + 5$ messengers and a symmetry which exchanges the $5(\bar{5})$ from the first set with $\bar{5}(5)$ from the second set. In this case the D-term arises at three loops.

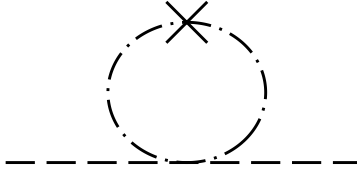


Figure 3: One loop epsilon-scalar counterterm graph contributing to the two loop soft scalar mass.

counterterm to correspond to minimal subtraction⁶.

Turning now to the soft scalar masses of the Standard Model fields, one finds that this counterterm contributes to the two loop soft scalar masses via a one loop counterterm graph shown in Fig. 3. This resulting graph is again logarithmically divergent (the factor ϵ^{-1} in the counterterm corresponding to (2.4) is cancelled after summing over the 2ϵ adjoint scalar multiplets running in the loop) and gives a contribution to the scalar mass of the form:

$$m_a^2 = -\frac{g_a^4}{128\pi^4} S_Q C_a \text{Str} M_{\text{mess}}^2 \log \frac{\Lambda_{UV}^2}{m_{IR}^2}, \quad (2.5)$$

where C_a is the quadratic Casimir ($(N^2 - 1)/2N$ for an $SU(N)$ fundamental; for $U(1)_Y$ C_a is Y^2 , with Y being the hypercharge of the Standard Model field involved). Λ_{UV} and m_{IR} refer to the ultraviolet and infrared cutoff respectively.

We now turn to considering the other graphs at two loops. These are identical to the graphs considered in [17] and are shown in Fig. 4. The contributions of these graphs can be calculated in a manner analogous to that in [17]. One finds that these graphs do not give contributions that are ultraviolet divergent. They do however give infra-red divergent contributions and these are cancelled by the infrared divergence in eq. (2.5). In fact it can be argued on general grounds of gauge invariance that no infra-red divergences can arise in the soft masses. The resulting cancellation between the different contributions is therefore a useful check on the calculation. Putting the contributions from the graphs in Fig. 3 and Fig. 4 together gives finally for the soft scalar masses⁷:

$$m_a^2 = \frac{g_a^4}{128\pi^4} m_f^2 C_a S_Q F(y_1, y_2, \Lambda_{UV}^2/m_f^2), \quad (2.6)$$

where C_a and S_Q again refer to the Casimir and Dynkin indices respectively as in eq.(2.5). The function F is given by:

$$F(y_1, y_2, \Lambda_{UV}^2/m_f^2) = - (2 y_1 + 2 y_2 - 4) \log \frac{\Lambda_{UV}^2}{m_f^2}$$

⁶We use the \overline{DR} [20] scheme where no “bare” mass for the epsilon scalars is introduced. We thank S. Martin for related discussions.

⁷The divergent contribution in eq. (2.6) has previously been obtained in refs. [15], [20].

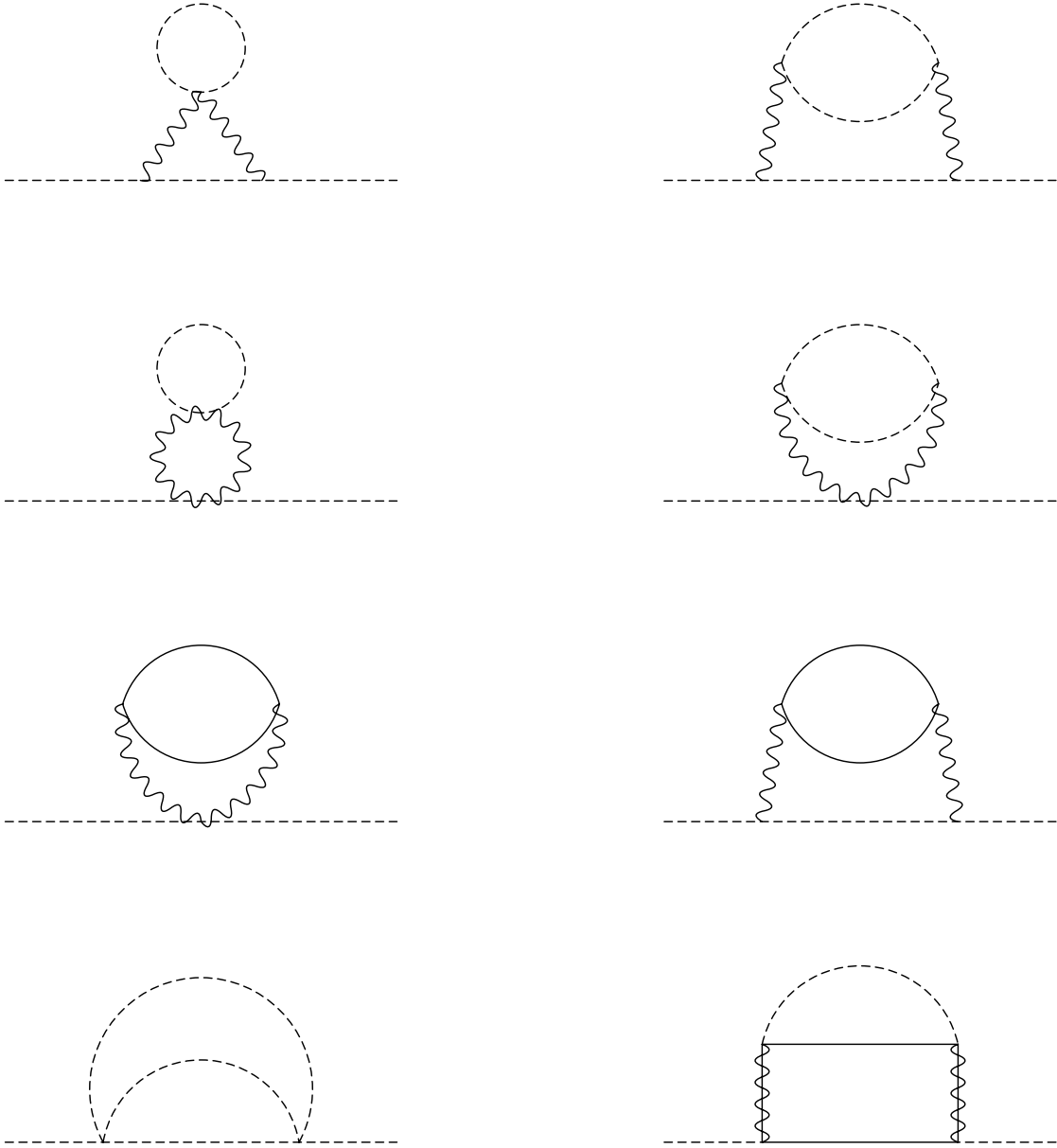


Figure 4: Two-loop messenger contributions to Standard Model soft scalar masses (the wavy lines denote both gauge bosons and epsilon-scalar propagators). The infrared divergence, present when $\text{Str } M_{mess}^2 \neq 0$, is cancelled by the counterterm graph of Fig. 3.

$$\begin{aligned}
& + 2 (2 y_1 + 2 y_2 - 4) + 4 x^2 (1 - x^2) (y_1 + y_2) \log y_1 \log y_2 \\
& + G(y_1, y_2) + G(y_2, y_1) ,
\end{aligned} \tag{2.7}$$

where

$$\begin{aligned}
G(y_1, y_2) & = 2 y_1 \log y_1 + (1 + y_1) \log^2 y_1 - 2 x^2 (1 - x^2) (y_1 + y_2) \log^2 y_1 \\
& + 2 (1 - y_1) \text{Li}_2(1 - \frac{1}{y_1}) + 2 (1 + y_1) \text{Li}_2(1 - y_1) \\
& - 4 y_1 x^2 (1 - x^2) \text{Li}_2(1 - \frac{y_1}{y_2}) .
\end{aligned} \tag{2.8}$$

As in the discussion of the gaugino masses, $y_1 = m_1^2/m_f^2$ and $y_2 = m_2^2/m_f^2$. $\text{Li}_2(x)$ above refers to the dilogarithm function and is defined by $\text{Li}_2(x) \equiv -\int_0^1 dz z^{-1} \log(1 - xz)$. It is easy to see that in the limit of vanishing supertrace and $x = -1/\sqrt{2}$ eqs. (2.3), (2.6) reproduce the results of refs. [17], [21].

As the first term in eq. (2.7) shows, in general for a non-vanishing supertrace of the messenger mass matrix, the soft scalar masses will depend on the ultraviolet cutoff. It is worth noting again that this ultraviolet divergent contribution arises from the one loop counterterm, eq. (2.5). Its presence indicates that in general the soft masses are sensitive to physics at scales higher than the scale of the typical mass of messengers ("the messenger scale"). This is to be contrasted with the case of vanishing weighted supertrace, when the typical momenta contributing to the scalar and gaugino masses are of order the messenger scale.

What the relevant ultraviolet cutoff is, will of course depend on the particular model. We will have more to say on this in the next section. Here we simply note that if there is a large hierarchy of scales in the supersymmetry breaking sector leading to the cutoff being much larger than the messenger scale, the term proportional to the supertrace in eq. (2.6) is the leading contribution to the soft scalar masses⁸. Then the general pattern one observes from eq. (2.6) is that the scalar mass squared is negative if the messenger supertrace is positive. Alternatively, the soft scalar mass squared is positive if the supertrace is negative. A discussion of the phenomenological relevance of this observation is also left for the following sections.

From the point of view of phenomenology, the main fact of importance is that the relation between the scalar and gaugino soft masses, characteristic of the minimal models of gauge mediated supersymmetry breaking [2], no longer holds in models with nonvanishing messenger supertrace.

⁸ At next to leading order in this case the logarithmically enhanced contributions arising at three loops could be comparable to the non-logarithmically enhanced contributions in eq. (2.6). We have not calculated these three loop contributions.

3 Ultraviolet sensitivity and phenomenological consequences.

In this section we address two main issues. Section 3.1 investigates the effects that cutoff the logarithmic divergence in the soft scalar masses, eq. (2.6), in the framework of both renormalizable and nonrenormalizable models of supersymmetry breaking. Section 3.2, addresses some of the phenomenological consequences of models of supersymmetry breaking with non-vanishing messenger supertrace.

3.1 What cuts off the logarithm?

3.1.1 Renormalizable models.

The ultraviolet divergence in the scalar soft masses indicates that these masses are sensitive to short distance physics and cannot be fully calculated within the low-energy effective theory. One needs to therefore go beyond the effective theory to the full underlying theory to estimate them. Roughly speaking one expects that if in the full theory there are additional fields that carry Standard Model quantum numbers and can play the role of heavy messengers, and if these heavy fields restore the full supertrace to zero, then the logarithmic divergence would be cut off by the scale of the heavy messengers (these heavy messenger fields could then also contribute to the masses through threshold effects). Whether this happens or not depends on the models under consideration and it is useful, in the discussion below, to distinguish between the case when the underlying theory is a renormalizable theory and when it is a nonrenormalizable theory. Examples of both types of theories of dynamical supersymmetry breaking exist in the literature.

We first consider the case when the underlying theory is renormalizable. In this case one can conclude that there must be extra heavy messenger fields in the full theory and, moreover, that the full weighted supertrace, after including the heavy messenger fields, must cancel. This follows from the following argument: If the full supertrace does not vanish the Standard Model soft masses will continue to be logarithmically divergent in the full theory and there would have to be a counterterm to absorb this divergence. Since the full theory is renormalizable such a counterterm would have to be renormalizable as well and would have to respect all the symmetries of the Lagrangian. Furthermore, this counterterm would involve a product of the Standard Model matter fields and fields from the supersymmetry breaking sector (the "spurion" fields). However, it is easy to see that no such renormalizable term can exist—since the soft masses are nonholomorphic, they must come from a term in the Kähler potential, and so the counterterm must necessarily have dimension greater than 4 [6]. On adding the contribution of the heavy messengers of mass m_H , the $\log(\Lambda_{UV}^2/m_f^2)$ term in eq. (2.6) will be

replaced by $\log(m_H^2/m_I^2)$.

In addition there could be threshold effects coming from these heavy messengers as well. These contributions to the soft scalar mass squares are proportional to $(\delta m_H^2/m_H)^2$, where δm_H is the typical splitting of the heavy supermultiplets. Whenever the ratio of the mass squared splitting, δm_H^2 , of the heavy supermultiplets to their mass, m_H , is of the same order as the corresponding ratio for the light messenger supermultiplets, the finite contributions of the heavy messengers will be comparable to those of the light messengers (the finite contribution of the heavy messengers can be additionally enhanced by their multiplicity [9], [10]). Whether or not such contributions are present is a rather model dependent question⁹. Generally, however, since the mass splitting of the heavy supermultiplets are not expected to be greater than the supersymmetry breaking scale, $\delta m_H \leq M_{SUSY}$, one expects that as the mass m_H increases, the finite contribution $(\delta m_H^2/m_H)^2$ of the heavy messengers becomes negligible.

3.1.2 Nonrenormalizable models.

We now turn to discussing nonrenormalizable supersymmetry breaking sectors. These typically contain interactions suppressed by some scale M_{UV} , and are themselves effective field theories valid below that scale. For the nonrenormalizable model to be a useful starting point in calculating the vacuum expectation values involved we need that both s and F_s are $\ll M_{UV}$. Here s and F_s are the vevs of the spurion field eq. (1.1) and represent the typical expectation values in the supersymmetry breaking sector.

Eqs. (1.1), (1.2), and (1.3) show that one expects the leading order contributions to the supertrace to be $\sim (F_s/s)^2$. In addition there could be subleading contributions which go like $(F_s/M_{UV})^2$. We now argue that the leading order contribution $\sim (F_s/s)^2$ must vanish when we include all the fields in the theory below the scale M_{UV} . In the previous section we used renormalizability to argue for the absence of a possible counterterm and therefore for the vanishing of the supertrace. This argument is not directly applicable here, since we begin with a nonrenormalizable model. Note though, that while nonrenormalizable counterterms might be allowed, they must still be polynomial in momenta and masses (the latter in the present context are dynamically generated). But a little thought shows that there is no polynomial counterterm¹⁰ that can account for the log-divergent contribution to the soft masses of the form (2.6) when $\text{Str } M^2 \sim (F_s/s)^2$. Hence, we conclude that in nonrenormalizable models, the leading logarithmic divergence is cutoff by heavy messenger fields with a mass m_H , $\Lambda \leq$

⁹These contributions are present e.g. in the $SU(N) \times SU(N-M)$ models (renormalizable for $N=4, M=1$ and $N=5, M=2$), where the mass squared splitting of the light messenger fields is much smaller than the supersymmetry breaking scale, and the ratio $(\delta m^2/m)^2$, is the same for light and heavy messenger fields.

¹⁰Nonpolynomial counterterms can be written—e.g. the D-term $\Phi^\dagger \Phi |\log S|^2$, where Φ is a Standard Model field and S a supersymmetry breaking sector “spurion” field with expectation value (1.1).

$m_H < M$ (Λ here is the scale above which the sigma model that describes the light messengers breaks down).

Besides the leading contribution to the supertrace though there can be, as was mentioned above, subleading contributions, proportional to $(F_s/M_{UV})^2$. These do not have to vanish—the corresponding polynomial counterterm is of the form $\Phi^\dagger\Phi S^\dagger S/M_{UV}^2$. The importance of these subleading contributions to the soft scalar mass in any model will depend on the ratio of the scales s/M_{UV} , the magnitude of the logarithmic enhancement, and the coefficients of the counterterms mentioned above.

3.2 Phenomenological consequences.

As noted at the end of Section 2, the logarithmically enhanced term in eq. (2.6) changes the relation between scalar and gaugino soft mass parameters, $m_{gaugino} \sim m_{scalar}$, typical of models with vanishing supertrace. The logarithmic contribution to the scalar soft masses is expected to dominate over the finite contribution—the finite threshold corrections due to both heavy and light messenger fields—in case there is a sufficiently large hierarchy of scales in the supersymmetry breaking sector. As we now discuss, this in fact puts a significant constraint on the class of viable models. For, as eq.(2.6) shows, when the supertrace for the light messengers is positive the logarithmically enhanced term provides a negative contribution to the scalar mass squares. Therefore when this term dominates the scalars are driven to acquire vacuum expectation values and, in particular, $SU(3)_c \times U(1)$ is broken—a clearly unacceptable outcome¹¹.

One obvious way to try and avoid this possibility would be to construct models where the supertrace of the light messengers is negative. The models with supersymmetry breaking-cum-messenger sectors that have been studied in detail so far, in particular the models of [9], [10], have all yielded a positive supertrace of the light messengers¹². This poses a serious problem for these models. In [9], for example, obtaining positive scalar mass squares requires, for $11 \leq N \leq 27$, the scale m_H to be 2.4 – 2.5 times the light messenger scale. This is clearly an unsatisfactory situation in which case the weak coupling analysis of the ground state is not even valid. We are however not aware of any general argument that requires the positivity of the light messenger supertrace¹³. This possibility might even be realized in a more exhaustive

¹¹Similar observations were made in [10].

¹²We note that L. Randall has recently constructed some models that have a negative supertrace of the messengers [12].

¹³It is easy to construct nonlinear sigma models that incorporate supersymmetry breaking and light messengers, in which the sign and magnitude of the supertrace is a free parameter (for example, in the $SU(N) \times SU(N-2)$ models this can be achieved by adding additional terms, allowed by all symmetries, to the Kähler potential of ref. [9]). However, we are not aware of any dynamical models to which these sigma models are a consistent low-energy approximation.

study of other vacua of known models. Finally, it is worth mentioning that even in models with the required sign for the supertrace, the logarithmically enhanced term might still pose a problem by driving the scalars much heavier and thus resulting in light gluinos.

In fact the logarithmic term in eq. (2.6) can be quite significant even when the ratio of the heavy to light messenger scales is not very large. To illustrate this consider an example consisting of two sets of messengers (both, say, in the fundamental representation of the relevant group). Further let the light messenger fields have a non-vanishing supertrace which is cancelled by the supertrace of the heavy fields thereby setting the full supertrace to zero. On choosing the ratio of the heavy to the light messenger fermion masses to be ~ 3 and choosing the light supertrace to be positive \simeq the light fermion mass in magnitude, one finds that the Standard Model scalar squared masses are generally negative. Further, if the sign of the light supertrace is reversed to be negative, the Standard Model masses now become generally positive. The logarithmic term can thus have a significant effect even for a small separation of scales between the heavy and light messengers and obtaining positive soft scalar masses in its presence is a significant constraint.

We conclude this section by briefly commenting on the "hybrid" models of [9]. In these models the soft scalar masses get comparable contributions from both gauge and gravitational effects. The negative contributions to scalar masses (arising from a positive light supertrace) from gauge mediation are then not necessarily a problem since they could be compensated by positive supergravity contributions. In fact they could lead to squarks and sleptons being lighter than gluinos—a novel and quite distinct spectroscopy¹⁴. For example, one can check that in the $SU(17) \times SU(15)$ models the (positive) supergravity contribution to the soft masses (due to the term that cancels the cosmological constant and to higher dimensional terms in the Kähler potential) is comparable to the (negative) log-enhanced contribution of the light messengers' supertrace. A fortuitous cancellation between these two contributions could then lead to squarks which are generically lighter than the gauginos. It would be interesting to study the renormalization group effects in these models in some detail—superparticle spectra with squarks much lighter than gauginos can not arise in (minimal) supergravity models.

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