New Mechanisms of Gauge-Mediated Supersymmetry Breaking

Lisa Randall

Center for Theoretical Physics
Laboratory of Nuclear Science and Department of Physics
Massachusetts Institute of Technology
Cambridge, MA 02139, USA
lisa@ctptop.mit.edu

Abstract

We introduce new mechanisms for the communication of supersymmetry breaking via gauge interactions. These models do not require complicated dynamics to induce a nonvanishing $F$ term for a singlet. The first class of models communicates supersymmetry breaking to the visible sector through a “mediator” field that transforms under both a messenger gauge group of the dynamical supersymmetry breaking sector and the standard model gauge group. This model has distinctive phenomenology; in particular, the scalar superpartners should be heavier by at least an order of magnitude than the gaugino superpartners. The second class of models has phenomenology more similar to the “standard” messenger sectors. A singlet is incorporated, but the model does not require complicated mechanisms to generate a singlet $F$ term. The role of the singlet is to couple fields from the dynamical symmetry breaking sector to fields transforming under the standard model gauge group. We also mention a potential solution to the $\mu$ problem.

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1 Introduction

There are essentially two ways to communicate supersymmetry breaking. Supergravity mediated supersymmetry breaking has the virtue of simplicity. Supersymmetry breaking can occur in a so-called hidden sector, and no contortions are required to transfer the breaking of supersymmetry to the visible sector because it is automatically accomplished via Planck-suppressed operators.

Gauge-mediated supersymmetry breaking on the other hand, conceivably possesses several advantages. Probably the most important from a phenomenological perspective is that flavor changing neutral currents are naturally suppressed. It is perhaps also somewhat more comfortable to have the physics of supersymmetry at lower energy scales, although the choice is nature’s and not ours.

So far however, gauge-mediated supersymmetry breaking appears not to be as attractive an option as one would like. Although models for communicating supersymmetry breaking exist, they are quite cumbersome. Furthermore, they generally have potentially dangerous color breaking vacua. From our perspective, the first problem is the more serious complaint, since it is hard to believe that nature has chosen the complicated mechanisms which are currently discussed in the literature.

There are several reasons why communicating supersymmetry breaking via perturbative interactions at a low-energy scale appears to require complicated structure. It is commonly accepted that the ideal scenario would somehow embed the standard model gauge group into a dynamical supersymmetry breaking sector in such a way that standard model particle superpartners automatically have the requisite supersymmetry breaking mass. However, the generation of such a model has proved difficult. The primary problem is that to obtain all the required soft supersymmetry breaking masses (including the gluino), one generally runs into a problem with a low-energy Landau pole due to the large number of flavors carrying standard model gauge charge.

For this reason, the idea of a messenger sector was introduced, which less directly communicates the breaking of supersymmetry to the visible sector. A messenger sector includes a vector representation of the standard model. The most popular of these models couples this vector representation to a singlet which has a nonzero $A$ and $F$ component, thereby transmitting supersymmetry breaking to the messengers, and consequently, to the visible sector. For example, a gluino mass is generated by the diagram of Figure 1, which clearly requires both an $A$ and $F$ type VEV in order to flip chirality on both the fermion and scalar lines.

*Even gravity-mediated supersymmetry breaking is complicated by the necessity for a singlet $F$-term to give a tree-level gaugino mass.*
This is readily seen explicitly, or by considering the \( \text{U}(1) \) carried by the vector quarks.

There are two beneficial features of a messenger sector, in addition to the eradication of the dangerous Landau pole. First is that the gluino diagram is generated at one loop, while the squark mass squared is generated at two loops, so that the resulting squark and gluino mass are of the same order of magnitude. Furthermore, the gaugino mass requires a single insertion of a supersymmetry-breaking \( F \)-term, while the squark mass squared requires two. If there is a singlet VEV, \( S \), which generates the messenger fermion mass, the gaugino mass is of order \( F/S \) and the squark mass squared is of order \( (F/S)^2 \), again a desirable relation. Together, these relations guarantee that the squark and gluino masses are about the same. The second advantage to models in the literature is that an explicit computation of the relevant Feynman diagrams \([4, 9, 10]\) shows that the squark mass squared arising from the single \( F \) term in the hidden sector is positive.

Despite these advantages, one might nonetheless reserve enthusiasm for these models, primarily because of the complications which are employed to give the singlet an \( F \)-type VEV. Singlets generally do not play a role in dynamical symmetry breaking models. A complicated scenario is generally required to communicate supersymmetry breaking to the singlet coupled to the messenger sector. When a fundamental singlet is coupled directly into the dynamical symmetry breaking sector, it appears to be nontrivial to couple the singlet to messenger quarks without introducing a flat direction in which squarks get a VEV. It is the generation of the \( F \)-type insertion for the singlet that seems to be the key problem for generating gauge-mediated models of supersymmetry breaking.

A nice alternative to fundamental singlets was proposed by Poppitz and Trivedi \([7]\), in \footnote{I will refer to squarks and gluinos explicitly; corresponding results apply to sleptons and charginos, etc.}
which the messenger sector is embedded directly into the supersymmetry breaking sector. The usual problem with the Landau pole is avoided because of the existence of two scales. Although many states carry standard model quantum numbers, many are heavy. The low energy sigma model can be chosen to avoid a bad Landau pole.

Although it is more compelling to have a model in which singlets are not introduced “by hand”, these models are in practice problematical. The first problem is that it is in fact nontrivial, though not impossible [11] to embed a group as large as the standard model as a global symmetry group into the dynamical symmetry breaking sector. The second problem is that in models with a hierarchy of scales, there is necessarily a mass range for which \( \text{Str}(M^2) \neq 0 \) where the supertrace is taken over the messenger fields. Explicit calculation [9, 10] for existing models in which \( \text{Str}(M^2) > 0 \) shows that this scenario generally implies negative mass squared for the squarks, unless parameters and models are carefully chosen.

My goal in this paper is to explore alternatives for communicating supersymmetry-breaking via gauge interactions which do not require an \( F \)-term for a singlet which couples directly to messenger quarks. In the models I present, there is a different paradigm for simplicity than the first one suggested, in which the gauge-mediated model embeds the standard model within the supersymmetry breaking sector. The paradigm is more similar to that suggested by hidden sector models, in which supersymmetry breaking can be communicated to other sectors in a fairly generic fashion, without structure which relies on a particular hidden sector model. In this paper we try to see how far we can get with simpler structure to the messenger sector and suggest examples. The major distinguishing feature of our models is that we permit there to be tree-level mass terms in the superpotential. Naively, this might seem to be counter to the philosophy of dynamical models, in which one avoids introducing mass scales by hand. However, we have learned in recent years that generating mass terms dynamically is very straightforward. Simple examples arise from compositeness, or dimension three Yukawa couplings matching onto mass terms due to strong dynamics [12]. A mass term could even be generated more prosaically from a Yukawa coupling to a field which obtains a mass because a mass squared scales negative. The only real requirement we would like to impose is that it is not necessary to take two different mass scales with entirely separate origin to be the same. This would amount to fine-tuning, and is the reason we believe the \( \mu \)-term requires some further explanation. In our theories, there will be qualitative requirements on mass parameters (that they be large or small compared to other masses) but different mass scales are not required to coincide.

The first class of models, discussed in Section 2, requires a messenger gauge group and “mediator” quarks which transform both under the messenger and standard model gauge groups. The squark and gluino masses both arise at high loop order, in such a way that the
squark is generically predicted to be heavier than the gluino by at least an order of magnitude. This would imply a relatively light gaugino (or heavy scalar) spectrum, subject to experimental verification. The mediator models can employ dynamical models of a messenger sector, in a way in which both problems mentioned above are solved.

In the second class of models, discussed in Section 3, I incorporate a heavy singlet “intermediary” field. The main distinguishing feature of this class of theories from other gauge-mediated models with singlets is the fact that the singlet does not acquire an $F$ term via complicated interactions and no messenger gauge group is necessary. The singlet is present in order to generate a higher dimension operator which connects the symmetry breaking sector to the messenger sector.

We briefly discuss phenomenology in Section 4. We conclude in the final section. An Appendix gives examples of dynamical supersymmetry breaking sectors which can be used in mediator models.

## 2 Mediator Models

The first class of models is based on the fact that a gluino mass can be generated by the diagram of Figure 2, instead of that of Figure 1\footnote{A similar diagram was considered in a somewhat different context in Ref. \cite{12}.}. (Of course one must also include the supersymmetric analogs; we present explicitly only the diagram in which supersymmetry breaking is communicated.) The necessary fields and interactions for such a diagram to exist are the following.

There is a weakly gauged global symmetry $G_m$ acting on fields in the dynamical supersymmetry breaking sector. This gauge group may or may not be broken, but a supersymmetry
breaking gaugino mass, $M_m$, for the gauge bosons of this group is essential. The existence of the messenger gauge group, and the existence of a supersymmetry breaking messenger gaugino mass, are the aspects of the model which depend on the dynamical symmetry breaking sector. Examples of models with these properties are discussed in an Appendix. Second, there are “mediator fields”, which we call $T$ and $\bar{T}$, in a vector representation of the standard model, which transform both under the messenger gauge group, $G_m$, and the standard model gauge group, $G_{SM}$ (or an extension thereof) (and therefore “talk” to both the dynamical supersymmetry breaking and visible sectors). In order to keep the number of flavors of the standard model small, we will restrict our attention here to $G_m$ having rank up to four, SU(2) being probably the simplest group to accommodate. For example, $T$ can transform under $SU(2)_m \times SU(5)_{GUT}$ as a $(2,5)$ and $\bar{T}$ as a $(2,\bar{5})$. Note that grand unification is not essential and the $T$’s can transform under the standard model gauge group, $SU(3) \times SU(2) \times U(1)$ instead. Third, there is a supersymmetric mass term $M_T T \bar{T}$. This might seem counter to the philosophy of dynamical models, in which mass scales are not specified by hand. The real requirement however is that the mass scales should not be introduced in a fine-tuned fashion. We will see shortly that $M_T$ is constrained to lie between $M_m$, the messenger gaugino mass, and the mass of the heaviest of the hidden sector scalars, $M_{DSB}$. If there is no large hierarchy then $M_T$ must accidentally agree with the scales in the supersymmetry breaking sector. If there is a large hierarchy between the above scales, as exists in some calculable models of supersymmetry breaking, or when supersymmetry is broken through higher dimensional operators, there can be a substantial range for $M_T$ and the model is natural.

With these ingredients, supersymmetry breaking can be communicated to the visible sector via the $T$ field so long as there can be a supersymmetry breaking $G_m$ gaugino mass. Examples where the messenger gaugino will be massive are discussed in an Appendix. In this model, the gluino and squark mass are both generated at high loop level. It should be noted that the complicated Feynman diagrams do not make the theory more complicated; they are present whether or not we compute them. However, because these are higher loop diagrams, we estimate, rather than compute the masses of the superpartners of the visible sector. This is sufficient to allow us to identify the dependence on couplings, loop factors, logarithms, and power dependence on masses. For this estimate, it is most expedient to divide the analysis into three possible ranges of parameters according to the relative sizes of $M_T$ and $M_{DSB}$, where $M_{DSB}$ is the mass of the heaviest field with a supersymmetry breaking scalar mass.

$M_T \sim M_{DSB}$: This mass range is not necessarily natural, as $M_T$ and $M_{DSB}$ have separate origin, counter to the philosophy espoused in the introduction. We include it for three reasons. First, pedagogically, it is simplest to first count loop factors, independently of mass suppressions or enhancements, which can be done most simply when all masses are compa-
rable. Second, it is a logical possibility which merits consideration. Third, even though we do not know the solution to the $\mu$ problem, there must be one; whatever mechanism relates the scale $\mu$ to the scale of soft supersymmetry breaking can in principle relate the $T$ mass to $M_{DSB}$.

The gluino mass is generated by the Feynman diagram of Figure 2. The gluino mass is nonzero, even without the presence of an $F$-term coupled directly to the messenger squarks $T$ and $\tilde{T}$. This is because the two-loop subdiagram which has external $\tilde{T}$ and $\tilde{\tilde{T}}$ states plays the role of $F_S$ in completing the Feynman diagram. The gaugino mass evaluates to a number of order $\alpha_s\alpha_m^2 F_{DSB}/(4\pi)^3 M_{DSB}$ as it occurs at three loops. Here we have used the fact that $G_m$ is a weakly coupled gauge group which does not play an essential role in the supersymmetry breaking dynamics. Therefore, $M_m$ (really the innermost loop of Figure 2 which is a self-energy diagram for the $G_m$ gauge boson) is loop-suppressed, and is of order $\alpha_m F_{DSB}/4\pi M_{DSB}$ (for this parameter regime there is not necessarily a distinction between these two mass scales) and is generated in the usual way the gluino mass is generated in visible sector models.

The squark mass squared is generated by four-loop diagrams not explicitly shown. Some of these four-loop diagrams can be readily identified by inserting a two-loop diagram of the standard sort with intermediate DSB states $[4, 9]$ that generates a supersymmetry breaking $\tilde{T}$ or $\tilde{\tilde{T}}$ mass into a two-loop diagram with intermediate $T$ states that generates the squark mass squared. The squark mass evaluates to a number of order $\alpha_s\alpha_m F_{DSB}/(4\pi)^2 M_{DSB}$. The sign of the mass squared is not known without an explicit calculation, but we expect one can choose parameters for which it is positive.

There also exist contributions to the $T$ squark mass which arise by inserting the effective "$F$" term which is generated at two loops. However, these contributions give rise to a six-loop contribution to the squark mass squared and are therefore negligible. This can be seen also from the form of the $T$ squark mass matrix. In addition to the diagonal supersymmetry breaking contribution to the $\delta M_\tilde{T}^2\tilde{T}\tilde{T}$ mass there is the off-diagonal $m_{LR}^2\tilde{T}\tilde{T}$ type mass. These mass squared parameters are of the same order of magnitude. However, $m_{LR}^2$ appears squared, whereas $\delta M_\tilde{T}^2$ appears once in generating a squark mass squared. Therefore the off-diagonal contribution is negligible and can be neglected.

We conclude that the ratio of gluino to squark mass is suppressed by $\alpha_m/(4\pi)$ in this model. The question then is what is a reasonable range for $\alpha_m$. The best situation is if $\alpha_m$ is as large as possible so that one does not run into naturalness problems. For $\alpha_m$ about 1, in which case $G_m$ can still be considered weakly coupled, this ratio is about 10, which is probably acceptable, and points to interesting predictions. If $G_m$ is not asymptotically free, a dangerous Landau pole will develop if $g_m$ is much bigger than 1. On the other hand, $G_m$ can have rank

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5I thank Bogdan Dobrescu for stressing the viability of large $\alpha_m$. 

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up to three or four, without running into problems with the running of SU(5). Even for SU(2), with six flavors of doublets (the minimum content required to have five flavors of messengers plus a doublet flavor in the DSB sector), the coupling does not run at leading order. So it is reasonable to expect that the messenger group can allow fairly big values of $\alpha_m$.\footnote{The other potential problem with large $\alpha_m$ is that the analysis of the supersymmetry breaking vacuum might need to incorporate the extra gauge group, and the separation between the DSB and messenger sectors would not be as clean. However the model would still break supersymmetry, since the equations of motion would still be inconsistent. Whether supersymmetry is broken should be independent of the ratio of couplings in the theory; if supersymmetry is broken at weak coupling it should also be broken when the coupling is somewhat bigger.}

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We conclude that there is a hierarchy between the gaugino and scalar masses (of like charges) of order 10. This prediction will persist for the other acceptable parameter range $M_T < M_{DSB}$ as we will see shortly. The precise ratio of scalar to gaugino mass depends on numbers expected to be of order unity which we have not incorporated. It is almost certainly a prediction of our model that the gauginos will be light and the scalars relatively heavy. We briefly discuss existing bounds in Section 4. To determine the naturalness of this parameter range also requires assumptions about the $\mu$ term and a more detailed calculation of the relevant Feynman diagrams. We interpret the large mass ratio as a prediction of our class of theories to be tested at future colliders. A more detailed analysis of the viable parameter range would be very worthwhile.

$M_T > M_{DSB}$: We next consider the possibility that $M_T$ is the heaviest mass scale. We now show that this possibility is not viable, because the gaugino mass would be suppressed by mass as well as loop factors in comparison to the squark mass, and is therefore too small.

This can be seen by an operator analysis, by explicit evaluation of the Feynman diagrams, or by an effective theory calculation of the Feynman diagrams.

Let us first consider the operator analysis. In terms of a spurion whose $F$-component breaks supersymmetry ($\Phi$) and which we also assume to have a nonvanishing $A$-component, we can construct the operators:

\[
\int d^4\theta Q^c Q^\dagger \Phi^\dagger \Phi \frac{M_T^2}{M_T^2} \tag{1}
\]

which gives rise to the squark mass squared and the operator

\[
\int d^2\theta W_\alpha W^\alpha \Phi \Phi \frac{M_T^2}{M_T^2} \tag{2}
\]

\footnote{$G_m$ should not be a $U(1)$ gauge group in any case, because the presence of the $T$ fields which transform under $U(1)$ and $U(1)_Y$ can generate dangerous kinetic energy mixing terms \cite{footnote}.}

\footnote{I am grateful to Ann Nelson for discussions.}
which contributes to the gluino mass. The two factors of $\Phi$ are required because we need both an $A$ and $F$-type insertion of $\Phi$ in the relevant Feynman-diagram. From the above, it is readily concluded that the squark mass goes like $F/M_T$ where as the gaugino mass is suppressed by $F/(M_T)^2$. The additional suppression by $M_T$ is undesirable since the gaugino mass is already light. In fact, there are larger contributions to the squark mass squared, not at all suppressed by $M_T$, so the ratio is even worse than is suggested by the operator analysis. This is understood from Feynman diagrams directly, or by including operators involving a nonvanishing $Str$. We briefly discuss the Feynman diagrams here, and reserve the operator analysis with nonvanishing $Str$ to the third case, where it is of more phenomenological relevance.

The $M_T$ suppression of the gluino mass can be readily understood from the Feynman diagram directly. The innermost loop can be thought of as generating a gaugino mass $M_m$ if the momenta running in the loop are less than $M_{DSB}$. However, because the relevant momenta from the full diagram are of order $M_T$, the result is suppressed by $F_{DSB}M_{DSB}/M^2_T$, where the first two factors were necessary to complete the innermost loop of Figure 2. This can also be seen from the effective theory below the $T$ mass scale. There is a two-loop diagram which generates an operator suppressed by two powers of $M_T$. Closing the loop of the light DSB fields again gives the suppression factor we have discussed.

One can also study the Feynman diagrams contributing to the squark mass squared. In fact one finds that the operator analysis above misses large contributions to the squark mass which are not mass suppressed at all. These can be understood as the $T$ scalar having an unsuppressed contribution to its mass splitting which does not decouple when computing the squark mass squared. (Further discussion of such effects are in the following section.)

We conclude that this case is not interesting from the point of view of gauge-mediated models.

$M_m < M_T < M_{DSB}$ For this range to exist requires that there is a hierarchy between the scales $\sqrt{F_{DSB}}$ and the scale $M_{DSB}$, the scale of mass for the heaviest multiplet in the DSB sector with a nonsupersymmetric realization. Many calculable models of supersymmetry breaking have such a hierarchy, as do models in which supersymmetry breaking occurs through the presence of nonrenormalizable operators in the superpotential. This is perhaps the most natural regime for our models.

The first point to understand is that the ratio of gluino to squark mass will not be suppressed by positive powers of $M_T$. Naively, this might have seemed to be the case because the gluino mass requires nonzero $M_T$ whereas the squark mass does not.

We can analyze the gluino mass in this model most simply, because we first integrate out the massive hidden sector fields (those with mass greater than $M_T$). This simply gives the
standard one-loop messenger gauge boson mass $M_m$ which can be considered “hard” below the scale $M_{HS}$. So the diagram contributing to the gluino mass is simply two-loop, as in Figure 3. The first inner loop generates a mass term of the form $m^2_{LR} \tilde{T} \tilde{T}$, where $m^2_{LR} \propto M_m M_T$ is generated by the inner loop. When inserted into the final loop, which is infrared convergent, the factor of $M_T$ from $m^2_{LR}$ is cancelled by an $M_T$ from the remaining loop (really one in the numerator divided by two in the denominator) to give a result which is independent of $M_T$ and depends on the same mass scale as the squark with no further mass suppression factors. This results holds only insofar as $M_T > M_m$. For $M_T$ smaller than $M_m$, one obtains the necessary mass suppression factor which yields a vanishing result when $M_T \to 0$.

The two-loop diagram which gives the gluino mass (once $M_m$ is considered “hard”) has the interesting feature that there is a logarithmic enhancement of the result due to the divergent inner loop. The result for the gluino mass is therefore of the order of $\alpha_s^2 \log(M_T/M_{DSB})^2 F_{DSB}/(4\pi)^3 M_{DSB}$.

The four-loop calculation of the squark mass is more difficult and subtle. Here it should be borne in mind that there are two scales of mass for hidden sector fields. For example, in the model of Poppitz and Trivedi, briefly discussed in an Appendix, the mass of the light fields in the low-energy sigma model is of order $\sqrt{F_{DSB}}$, whereas there are also heavy fields with mass scale set by $M_{DSB}$ (times a gauge or Yukawa coupling). The heavy fields couple to light fields with nonvanishing $F$-components, and contribute to both gaugino and squark masses. However, the diagrams with internal light fields for DSB states (lighter than $M_T$) also contribute to the squark mass, without mass suppression! This seems to contradict the wisdom of the operator analysis which suggested that for $M_T > M_{DSB}$, the Feynman diagrams contributing to the squark mass squared would decouple. The essential difference is that the previous discussion only applies to $F$-type contributions to the mass squared. In models with a separation of scales, there is very likely also a contribution from nonvanishing $Str(M^2)$ between the two mass scales of the DSB scalars. This is in fact the case in the

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*I am grateful to Erich Poppitz for discussions, and for sharing the results of [10] before publication.*
models of Ref. [7] for example. And the $T$ sector always has nonvanishing $\text{Str}(M^2)$. The mass for the $T$ scalar contributes in turn to the squark mass squared, and for this mass range should be the dominant contribution since it is logarithmically enhanced. Before examining this contribution in more detail, let us first understand the nondecoupling of the $\text{Str}(M^2)$ term.

Before we argued that when $M_T$ is large, the squark mass squared will be suppressed by $F_{\text{DSB}}^\dagger F_{\text{DSB}}$. Because this is dimension four, the dimensions are compensated by two factors of $M_T$, yielding the mass suppressed results we found earlier. However, $\text{Str}(M^2)$ is dimension two. If $\text{Str}(M^2)$ has magnitude of order $F_{\text{DSB}}$ (as it does in Ref. [7]), it can be thought of as $F_{\text{DSB}}^\dagger F_{\text{DSB}}/|F_{\text{DSB}}|$, where the dimensionful factors in the denominator are already supplied, and are the lighter scale $|F_{\text{DSB}}|$. So whereas the contributions due to $F$-type insertions decouple when inserted on a heavy internal line, this is not true for the contributions proportional to $\text{Str}(M^2)$. Furthermore, the diagrams in which $F$-terms are not inserted on the heavy line do not decouple, as can be seen by routing the momentum on the two-loop diagram contributing to the $T$ mass so that there is no common loop momentum between the $T$ line and the line on which the $F$ terms are inserted. The fact that $\text{Str}$ insertions do not decouple when inserted on a heavy line whereas $F$ type insertions do can be readily seen from a one-loop computation (simpler than the two- or four-loop calculations relevant to this model). Consider a quartic interaction between two heavy fields (mass $M_H$) and two light fields, in which there are both supersymmetry breaking $F$-type mass for heavy fields $\tilde{H}$ and $\tilde{\bar{H}}$ and supersymmetry breaking $\delta M^2$ masses for $\tilde{H}$ and $\tilde{\bar{H}}$. One can calculate the mass squared for the light fields obtained by closing the heavy fields in a loop and inserting $F$ twice. There are three propagators and the result is suppressed by $M_H^2$. On the other hand, when one calculates the contribution from inserting $\delta M^2$, there are only two propagators—the result is not only unsuppressed; it is divergent, and generates a logarithmic contribution to the running of the scalar mass squared for the external field. These results are nothing new; they are the reason gauge-mediated models are insensitive to high scales in general and supergravity models are not. The distinction here is that there exists a nonvanishing $\text{Str}$ contribution from a gauge-mediated model as well in certain cases when there is a separation of scales between fields which both contribute to the dynamical symmetry breaking sector, as occurs for example in models with a dynamical messenger sector. And there is always a nonvanishing $\text{Str}$ over the $T$ sector.

Having established that both the heavy and light fields in the DSB sector contribute to the squark scalar mass squared, we see that the essential difference between the first estimate for the squark mass squared (when $M_T \sim M_{\text{DSB}}$) and the case we are considering here is that the result can be logarithmically enhanced, and furthermore that we can establish the sign for the scalar mass squared if it is indeed the logarithmically enhanced piece which dominates.
The dominant contribution in this case can be thought of as arising in two steps. First, the scalar $T$ fields get a supersymmetry violating mass squared at two loops. Notice that this mass is such that $\text{Str}(M_T)^2 \neq 0$. The $T$ scalar mass is then inserted into another two-loop diagram which generates the mass squared for the squarks. The final contribution can be obtained from the renormalization group equations

$$\frac{d\delta M_T^2}{d(\log(\mu/M_{DSB})^2)} = \frac{g_m^4}{256\pi^4} (\text{Str}(M^2))_{DSB} C$$  

$$\frac{d\delta m_{\tilde{q}}^2}{d(\log(\mu/M_{DSB})^2)} = \frac{g_s^4}{128\pi^4} \frac{4}{3} \delta M_T^2(\mu)$$

Here $C$ is a group theory factor. In Eq. 3, we have taken the group theory factor for $SU(3)$ explicitly, and have used the fact that $\text{Str}(M^2)$ for one doublet flavor of the $T$’s is $2M_T^2 + 2M_2^2 - 4M_T^2$, where $T_1$ and $T_2$ are the mass eigenstates. When integrating the above, we will be interested in the result when $\mu = M_T$ for which the logarithm is negative. The sign of $m_{\tilde{q}}^2$ depends on the sign of $\text{Str}(M^2)$ in the DSB sector. If there is a dynamical model with a separation of scales as the DSB sector and the $\text{Str}$ is positive over the scales between the two masses, the mass squared for the squarks and sleptons will be positive. Furthermore the mass squared is enhanced by the factor $\log(M_T/M_{DSB})^2$, as can be seen by integrating the renormalization group equations.

It is an important requirement for the mediator model in the parameter regime $M_m < M_T < M_{DSB}$ where the logarithmically enhanced contributions to the squark mass squared dominate that the dynamical symmetry breaking sector gives positive $\text{Str}$. The $T$ scalars are then less massive than their fermionic counterparts, yielding negative $\text{Str}(M_T^2)$, which in turn gives positive mass squared to the squarks. If the only contribution to the $T$ mass squared came from nonvanishing $F$ terms, the $T$ scalars would be more massive than the fermions and the squark mass squared would be negative.

It is essential that $\text{Str}M^2$ is not parametrically larger than $(F/S)$ from the light fields, since the latter determines the gluino mass, whereas the former determines the dominant contribution to the squark mass. This turns out to be true of the models of Ref. [7], and should be a generic feature of dynamical models with a similar separation of scales.

We find that dynamical messenger sectors work much more neatly in conjunction with a model like this one, which gives the squark mass squared in two stages. Were the messenger sector to give mass squared directly to the squarks, it would generically be negative, though this conclusion can perhaps be escaped in specific models.

It might also be considered an advantage of this conjunction of the $T$ fields with dynamical models that the global symmetry of the DSB sector does not need to be sufficiently large to incorporate $SU(3) \times SU(2) \times U(1)$; $G_m$ can be as small as $SU(2)$. This can sometimes permit
a lower scale of dynamical supersymmetry breaking, for which higher dimensional operators which yield soft scalar masses (which can be flavor changing) are suppressed.

The remaining ingredient to consider for this class of models is to establish the nonvanishing messenger gaugino mass. Since the messenger gaugino mass is getting a mass exactly as the gluino would in more direct models of gauge-mediated DSB, but without having to embed a group as large as the standard model, we know this is possible. However, we wish to emphasize the special role of dynamical models, and also to suggest that even models in which the messenger gauge group is broken might work. Discussions of two models can be found in an Appendix.

3 Intermediary Models

Having considered a model in which it is gauge and not superpotential interactions which communicate supersymmetry breaking, we now consider a class of theories in which messenger gauge interactions are not necessary. We furthermore do not require a complicated mechanism to generate a nonvanishing singlet $F$-term. The model does however incorporate singlets, or a dynamically generated dimension four operator in the superpotential or the Kahler potential. The potential drawback to this model is the existence of additional local minima which do not break supersymmetry or break standard model gauge groups. Although we can argue (like in models [4, 5, 6]) that these vacua are local minima, to make the desired vacuum correspond to the deepest minimum could require additional structure [8]. The advantage of the model is that the singlet field functions very simply to produce an operator which permits direct communication of supersymmetry breaking to messenger fields.

The field content of the model is as follows. First, we assume the existence of a model which breaks supersymmetry dynamically. In general, such models are chiral. However, we are interested in models which include at least one vector representation, $V, \bar{V}$, of the gauge group whose dynamics is responsible for breaking supersymmetry. The requirement is that $V\bar{V}$ has a nonvanishing $F$-term. This can probably happen in many models.

A specific example including a vector representation which is worked out in the literature is in Ref. [13]. This model has the field content of the original ADS model based on SO(10) with a singlet 16 [16], but has an additional H(10) included. Murayama assumed this model can be analyzed perturbatively (in reality there can be an unbroken strongly interacting gauge theory so the analysis is not completely reliable [17]) and analyzed the vacuum. According to his result, the invariant $H^2$ has a nonvanishing $F$-term for small mass. The theory could not be perturbatively analyzed with large mass.

Another example in which there are vector representations is the model of Refs. [18, 19].
Here the vacuum was not analyzed for generic parameters, so it is not clear that the minimum
has nonvanishing $F$ term for one of the SU(2) mesons, though it is possible.

In both of these examples, the vector field had nontrivial Yukawa couplings in the superpotential. We suggest that the phenomenon of nonvanishing $F$ components for a vector field might be much more generic. Suppose we take a model which breaks supersymmetry through a combination of a superpotential generated by gaugino condensation and perturbative terms which lift all flat directions. Add to this model a massive vector multiplet with the only tree-level superpotential coupling being $m_{\nu}V\bar{V}$, which clearly lifts all new potential flat directions. This model clearly still breaks supersymmetry, since the old equations of motion are still inconsistent. This can be seen most simply by first integrating out the massive $V\bar{V}$ field.

Now consider the full theory, including $V$ and $\bar{V}$. We expect that in general $(V\bar{V})_F$ has a nonzero expectation value. Notice that $V$ and $\bar{V}$ have nonvanishing expectation values due to the dynamical term, and the mass term removes any potential flat directions involving the $V$ fields. If $V$ and $\bar{V}$ are heavy, this $F$-term is mass suppressed. This can be seen by perturbing the theory around the vacuum with vanishing $F$ term for the $V$ fields. There is a mass-suppressed tadpole for the $V$-field, which induces a mass-suppressed $F$-term. The power dependence of the mass-suppression is model-dependent. If $V$ is light, the $V$ field constitutes an approximate flat direction. We expect all vacuum expectation values and $F$-terms to be governed by a combination of its mass and the dynamical scale. The nonzero $F$ term is an assumption which can be verified in a detailed analysis. A priori, there is no reason there should not be a comparable $F$ term for $V$ as for other fields in the dynamical symmetry breaking sector, if $V$ is light, and a mass-suppressed value if it is heavy. A mass suppressed $F$-term should still be adequate, so long as the mass is not too big and the $F$-term is nonzero.

It is not necessarily essential to have a mass term in models with vector representations. The first two examples we gave did not involve mass terms. What is important is that all the flat directions are lifted. This is accomplished through Yukawa interactions in the first two models mentioned. In the absence of a specific model, the mass term is more generic, but not necessarily essential.

To complete the model, and transmit supersymmetry breaking to the visible sector, we assume the existence of two singlet fields, $S$ and $\bar{S}$ and a vector-like messenger representation of the standard model or a GUT extension, $Q$ and $\bar{Q}$. The superpotential contains $S V \bar{V} + M_{S} S S + S Q \bar{Q} + m_{Q} Q \bar{Q}$. Having two singlet field increases the likelihood that dynamical interactions generate the necessary mass terms and couplings and furthermore prevents potentially dangerous interactions between the $V$ and $Q$ fields. Here we have not explained the absence of other terms permitted by symmetries, but assume this will be explained by a
more fundamental theory. The fact that we want all the couplings to be renormalizable (not suppressed by $M_{Pl}$) is the reason we introduced a vector-like representation into the dynamical symmetry breaking sector. The large singlet mass is necessary to separate the desired minimum in which supersymmetry is broken and field values are independent of $M_S$ from the undesirable minimum with large field values which grow with $M_S$. One might hope that in a dynamical supersymmetry breaking model which has a singlet as an integral part of the model that no additional singlet would be required. However, without an additional singlet, we have found there is a new flat direction which destroys the model as a candidate for the supersymmetry breaking sector. The mass for the $Q$ in the superpotential above is required to keep the desired vacuum stable.

The simplest way to analyze what happens is to first integrate out the massive scalars. In this effective theory, there is a dimension four operator in the superpotential, $V\bar{V}QQ/M_S$. If $V\bar{V}$ has a nonvanishing $F$ term, and $Q$ has a mass, this model will work identically to the usual gauge-mediated models. No complicated couplings are required in order to generate an $F$ term for the singlet. At this point it should be clear that another possible model just incorporates the dimension four operator directly. However, if the scale is less than $M_{Pl}$, some explanation is warranted. The operator could be generated by composite interactions, but there would be common “preons” to hidden and visible sector fields, and one would also expect there to exist fields which transform under the gauge groups of both sectors. This can be dangerous, so we have instead considered the model with singlets. An alternative possibility is that the necessary operator arises from a Kahler potential coupling involving the $Q$ and fields from a dynamical sector, which is suppressed by a dynamical scale. Again, we have yet to realize this possibility, and have therefore generated the interaction via singlets.

It is readily seen in the effective theory that the old vacuum is still a local minimum when $Q$ and $\bar{Q}$ vanish, since all the old equations of motion for fields in the dynamical symmetry breaking sector remain valid. However, in the absence of a mass term for $Q$, we see that the mass term $QQ$ is suppressed by one power of $M_S$, while the mass term for $QQ^\dagger$ is suppressed by two powers, so there would be an unstable field direction. This is readily eliminated if the mass term for $Q$ is sufficiently large. Our minimum is probably not the global minimum, which can occur for field values of order $M_S$. Since the existing models are based on local not global minima in any case we are no worse off in this regard.

One can also analyze the vacuum in the theory with the $S$ and $\bar{S}$ fields still present. One finds similar conclusions, except that there is a nonvanishing vacuum expectation value of $S$ at order $1/M_S^2$. This means that the vacuum we want is shifted slightly, with the vacuum expectation value of fields shifting at order $1/M_S$.

Here we have found a simple generic mechanism to transmit supersymmetry breaking
to the visible sector. We have avoided the complications required to generate \( F \)-terms for the singlet. It would be nice to explicitly verify the nonvanishing \( F \)-terms for the vector representation. It is not hard to find a theory with the necessary number of vectors which can be analyzed in a perturbative regime; the analysis is however numerically complicated.

4 Mass Scales and Phenomenology

The mediator models we have discussed have a distinctive mass spectrum which should be readily distinguishable from that of other models of gauge-mediated supersymmetry breaking. The intermediary models on the other hand give phenomenological signatures very similar to models which have already been considered. We briefly discuss the mass scales and consequences for our models here and leave a more detailed analysis to further investigation.

We first discuss the scale of supersymmetry breaking. This is relevant because it will determine whether there will be photonic decays inside the detector [20, 5, 21]. This scale also determines the importance of potentially flavor violating scalar masses which can arise due to higher dimensional operators in the Kahler potential, which can be significant when the dynamical supersymmetry breaking scale is high.

In all cases, one anticipates a fairly high scale of supersymmetry breaking in the mediator models. This is because the gaugino mass arises at three loops. If we constrain the gluino mass to be about 100 GeV, this would imply (with \( \alpha_m \sim 1 \)) a supersymmetry breaking scale at least of order \( 10^6 \) GeV, too high to be likely to permit photon signatures. In a light gluino scenario, this scale might be smaller by an order of magnitude. If the hidden sector is provided by a dynamical supersymmetry breaking sector with a hierarchy of mass scales, one expects an even higher scale of supersymmetry breaking. For example, in the model of Ref. [7] discussed in the Appendix, if \( N = 5 \), the hierarchy of mass scales must be provided by a small Yukawa coupling, \( \tilde{\lambda} \). The scale of supersymmetry breaking is then higher by \( 1/\sqrt{\tilde{\lambda}} \). If \( N = 7 \), the scale of supersymmetry breaking is substantially higher because the supersymmetry breaking communicated to the messengers is suppressed by a ratio of mass scales. It is readily seen that the scale of supersymmetry breaking can be as high as \( 10^9 - 10^{10} \) GeV (where we have included the logarithmic enhancement of the gaugino mass). This is quite high, but should be consistent in our models in which the scalar masses are an order of magnitude higher than gaugino masses (for like charges). The danger is that there can be tree level flavor changing contributions to the scalar masses; for the \( N = 7 \) model they are suppressed but interesting. Clearly, the scale of supersymmetry breaking is model dependent. In all cases, it is likely to be large, but the precise value depends strongly on the dynamical symmetry breaking sector. The hierarchy of mass scales present in the dynamical messenger sector models (here
the messengers are for the gauge group $G_m$) is not an essential ingredient. The hierarchy played two roles in Ref. [7]. First, it allowed for a calculable model and second, it allowed for many states with standard model gauge charge to be heavy. The calculability of the model is not necessarily essential and potential Landau poles are not a problem since it is $G_m$ and not $SU(3) \times SU(2) \times U(1)$ which is the gauged global symmetry. However, a hierarchy of mass scales between $M_{DSB}$ and $\sqrt{F_{DSB}}$ might be a desirable feature in our models because it permits a larger range for $M_T$. We therefore expect a high scale for $F_{DSB}$, sufficiently high that photon signatures will not occur, but sufficiently low that nonrenormalizable terms are sufficiently small that FCNC effects are not too big.

The scale of supersymmetry breaking in the intermediary models is determined by the unknown mass parameters. So although supersymmetry breaking is transmitted to the messengers at tree level, the scale of supersymmetry breaking is likely to be high. We conclude that photon signatures are probably not a signature for either type of model.

The best tests of the models are the mass spectrum. The mediator models imply that the gauginos which are superpartners to the standard model gauge bosons are the lightest of the superpartners (of like charge), with the scalars of corresponding charges at least an order of magnitude heavier.

A very light gluino scenario [22] might still be viable. However, since $M_1$ and $M_2$ are also very small, such a scenario is very constrained in light of LEP 1.5 (and LEP 2) results, on top of any direct bounds on the gluino itself. In exploring the phenomenology of the light gluino scenario of this model, it should be remembered that in addition to the one-loop contributions to the gaugino masses with intermediate standard model superpartners, there is also the three-loop contribution we have already discussed. If $\alpha_m$ is of order $1/(4\pi)$, these should be of comparable importance.

A less restrictive parameter range will occur if $\alpha_m$ is of order unity. In this case, one can have a phenomenologically acceptable spectrum with the gaugino masses near their current experimental limit. The upper bounds on masses arise from naturalness considerations which would need to be redone for the mediator models. The lower limits are determined by recent LEP 1.5 results and by the gluino mass limits. In models with a grand unified mediator spectrum, the gaugino masses will be related by the gauge couplings, and the limit on any gaugino will imply limits for the other gauginos as well. The strongest bound on a gluino which is lighter than the squarks comes from D0 [23] and is 144 GeV. This bound was derived however for particular parameter choices. The most recent CDF [24] publication gave results for various parameter choices; in some regions of parameters there were no limits. UA2 gave a bound of 79 GeV [25] which applied for photino mass less than 20 GeV, which will be the case for the simplest models. It is conceivable there exists a gluino window for heavier masses.
when one does not assume the spectrum of mediators is grand unified. Because this can lead to other complications (a nonvanishing $D$-term for $U(1)_Y$ term for example), it is probably reasonable to assume the UA2 bound applies. It would also be worthwhile to extend the D0 bound to a broader parameter range. The current gluino mass bound (assuming the light gaugino scenario is not viable) is presumably on the order of 100 GeV.

However, the stronger bound on parameters will come from LEP 1.5 and LEP 2 results which will exceed the gluino mass bound if GUT relations are assumed. The most recent published results \cite{20} indicate a bound on $M_2$ between about 20 and 30 GeV. Because the sneutrino and slepton would be heavy in this scenario, the bounds from chargino and neutralino constraints can be readily interpreted as a bound on the gaugino mass matrix parameters. The bound is strongly dependent on $\tan \beta$ and favors a small value. It also favors a negative and small value for $\mu$. The naturalness of the model should be analyzed taking into account the chargino and neutralino constraints which will bound the parameters $\mu$ and $M_2$ and favor smaller values of $\tan \beta$ to maintain small $M_2$ in order to keep the scalar spectrum not too heavy.

For a particular value of $\alpha_m$ and the gauge group $G_m$, there will be a relation between the scalar masses and the gaugino masses. This will however be model dependent as there are in general contributions not only from the $F$ terms in the dynamical supersymmetry breaking sector, but from $Str(M^2)$ in the dynamical supersymmetry breaking sector, whereas only the $F$ terms contribute to the gaugino mass. The qualitative prediction is that the matching conditions for the scalar mass is on the order of ten times bigger than in the standard phenomenological analysis of gauge-mediated models. A more comprehensive study of the viability and phenomenology of our parameter range would be worthwhile.

The spectrum of the intermediary models is more similar to that previously investigated, in Ref. \cite{21} for example. The mass of the messenger quarks is however a free parameter, unrelated to the dynamics of supersymmetry breaking and its communication to the messenger squarks via the singlet.

5 Generating $\mu$

We have not yet addressed the other major stumbling block to gauge-mediated supersymmetry breaking, the generation of a $\mu$ term. This is another place where models seem to be complicated. The solution to the $\mu$ problem can be the same as in previous models of dynamical supersymmetry breaking. We present another potential solution which could also apply to other models and is really tangential to the rest of this letter.

A possible mechanism for generating a $\mu$ term is the following. Introduce a singlet field $M$
and charged fields under an $SU(N_X)$ nonabelian gauge group $X(N_X)$ and $\bar{X}(\bar{N}_X)$. Assume the superpotential contains $M H_u H_d + M X \bar{X} + M^3$ where $H_u$ and $H_d$ are the standard model Higgs superfields. Also assume that $X$ and $\bar{X}$ have weak scale masses which were generated in our model through $T_X$ fields, which transform under the messenger gauge group and the group under which $X$ transforms. (In the mediator models, it would be best for the $T_X$ and $T$ masses arise from a common confinement scale so that the mass of the $X$ fields is also weak scale. This can happen for example if $T$ and $T_X$ are composites of common preons transforming under $SU(2)_m$ as well as preons carrying either standard model of $SU(N_X)$ gauge charge).

We expect the $M$ scalar mass to be zero at the matching scale (since it is gauge-neutral) but to obtain a mass squared upon renormalization group running (more quickly for larger $N_X$). This will scale the $M$ mass squared negative, and should lead to a VEV for $M$ of order the weak scale, as in no-scale scenarios.

The $\mu B$ term is generated because the $M^3$ term implies $\langle F_M \rangle \sim M^2$. If $\lambda \sim 1$, this is the correct relation between $\mu$ and $\mu B$. This model is similar to that presented in Ref. [4]. The problem for the simplest model there was that the superpotential, which only has dimension three operators, preserves an $R$-symmetry (this was the U(1) identified in Ref. [4]) and therefore contains an associated $R$-axion which couples to the $Z$. This problem is readily solved here because the $R$-symmetry is anomalous with respect to the gauge symmetry under which $X$ transforms, so the pseudoscalar can be raised above the $Z$ mass through instanton-effects.

6 Conclusions

The lack of elegance to visible sector models is the chief reason to view them with skepticism. In this paper we have explored the question of whether there can exist alternative mechanisms for communicating supersymmetry breaking. In the mediator models, there is no need for singlets coupled directly to messenger quarks or for a complicated superpotential, and furthermore, there should be no new color or charge breaking minima. It should be noted that we have assumed the $R$-axion intrinsic to visible sector models is given a suitably high mass from supergravity and cancellation of the cosmological constant [27]. No new superpotential couplings are required, aside from that which gives the $T$ and $\bar{T}$ superfields a mass. Furthermore, the global symmetry of the Lagrangian, which is easy to obtain in many models of dynamical supersymmetry breaking, is not necessarily preserved by the dynamical supersymmetry breaking vacuum. Gauge interactions and heavy fields which transform under a messenger and standard model gauge group suffice to communicate supersymmetry breaking. However, for this class of models, there was either a narrow range for $T$ mass or a dynamical
messenger gauge group was required. Furthermore, the mass range which we predict will
probably not produce as natural a Higgs sector as that considered in Ref. [21], but should be
explored, and is subject to experimental verification.

The intermediary models are much simpler as given. Supersymmetry breaking is directly
communicated via a dimension four operator in the superpotential. The massive singlet is
necessary only in order to generate this operator. There are several mass parameters but
they are fairly arbitrary, with the constraint being only that the supersymmetry breaking
scale should not be so high that gravity-mediated supersymmetry breaking is of comparable
importance. Nonetheless, it would be of interest to produce these mass terms dynamically,
and work is underway here. It would furthermore be of interest to verify the dynamical
assumptions which were made; in particular, it would be good to verify the nonvanishing $F$
terms for the vectors.

Clearly, there is much more work to be done, in particular a better understanding of
the solution to the $\mu$ problem, and a more thorough investigation of the phenomenology of
mediator models. Nonetheless, we believe it is quite interesting that there can be very different
scenarios for gauge-mediation of supersymmetry breaking.

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8 Appendix: Generating $M_m$

Generating $M_m$ is not hard. However, analyzing the vacuum in nonperturbative supersymme-
try breaking models is often stymied by the fact that the Kahler potential is not determined,
so the location of the vacuum is uncertain. For this reason, it is often difficult to ascertain
the global symmetry group which remains at the minimum of the potential. We will argue
that it is not essential that the weakly gauged $G_m$ remains unbroken, although it is simplest
to first consider this possibility. The supersymmetry breaking mass for $G_m$ is generated akin
to the mass of the true gauginos, through a diagram like that of Figure 1.

We take the model analyzed in [7] as our first example. This model might be more
complicated than necessary, but since it has already been completely analyzed, it is a simple
model for us to consider. Furthermore, it has the separation of scales which could lead to
the most natural models of the type considered in Section 2. The models considered in Ref. [7] employed an $SU(N) \times SU(N - 2)$ gauge group, with odd $N$ and were shown to break supersymmetry. The couplings in the superpotential can be chosen to preserve an $SP(N - 3)$. In order to incorporate the standard model gauge group, it was necessary to take $N \geq 11$. However, we only wish to incorporate the gauge group $G_m$ which we take to be SU(2) so we can take $N$ as small as 5, although to make the theory perturbative without a small Yukawa coupling prefers $N = 7$.

This model establishes that it is possible to generate supersymmetry breaking gaugino masses for $G_m$ with a reasonable choice for the supersymmetry breaking scale. The above model is special in that one can determine the low-energy vacuum and ascertain that a global symmetry is preserved. The soft supersymmetry breaking mass and $F$ type terms have been calculated, and it was established that for the light fields that the singlet VEV scales as a parameter $v$ (related to more fundamental scales of the theory), the $F$ term for the light fields scales as $v^2(v/M)^{(N-5)}$ (or with a Yukawa for $N = 5$), and that $Str(M^2)$ over the light fields scales as $v^2(v/M)^{(N-5)}$. Here $M$ is likely to be $M_{Pl}$ but could be some other scale. Therefore the mass splittings and $F$ type vevs contribute the same order of magnitude to the squark and gaugino mass, aside from the logarithmic enhancements discussed in the previous section.

The next model we consider was given by Affleck, Dine, and Seiberg [28] and by Dine, Nelson, and Shirman in Ref. [5] and analyzed by ter Veldhuis in Ref. [29]. The model has two 10’s and two $\bar{5}$’s and a tree-level superpotential $W = 10_1 \bar{5}_1 \bar{5}_2$ (5) where this is the most general superpotential allowed by the symmetries. This superpotential preserves an SU(2) global symmetry (as pointed out in [5]) which is spontaneously broken [29]. In addition to the supersymmetric mass for the SU(2) gauginos, there should also be a supersymmetry breaking mass. For example, both the $A$ and $F$ components of 10$_1$ should be nonzero at the minimum of the potential, so the diagram of Figure 1 should generate a supersymmetry breaking gaugino mass for the SU(2) gauginos. Because of the supersymmetric contribution to their mass, the result for Figure 3 will be different. However, no large ratio is expected for values of the Yukawa coupling of order unity, and furthermore, any ratio can be accommodated by adjusting the scale of supersymmetry breaking.

A model of this sort without a separation of scales will require the $T$ mass to be of the same order of magnitude as the mass of fields in the DSB sector. It is interesting nonetheless that a model with a broken global symmetry group can also work. The danger with such a model is that the $D$-terms can be nonvanishing. This can potentially induce a VEV for the $T$ or $\bar{T}$
field which would be a phenomenological disaster. This will not occur so long as $M_f^2 > g^2 D_m$, where $D_m$ in this case is the $D$-term for SU(2) at the supersymmetry breaking minimum. The more serious danger for a nonvanishing $D_m$ is that the scalar mass will be generated at tree level or at one loop, making the gaugino to scalar mass ratio completely unacceptable. Therefore we require vanishing (or loop suppressed) $D_m$. This can follow because the gauge group is preserved, or because there is a charge-conjugation or other symmetry which protects the $D_m$ term. In the particular model just considered, this requires that a custodial symmetry is preserved at some level, which requires a small coupling ratio $g$.

It should be noted that it is not always true that there is a supersymmetry breaking gaugino mass. For example, in the $SU(6) \times U(1) \times U(1)$ model of Ref. [6], the messenger U(1) gaugino is massless. However, in general, the U(1) gaugino should pick up a supersymmetry breaking mass. A simple example of this is the $SU(3) \times SU(2) \times U(1)$ model of Ref. [4, 5]. However in this case there is a dangerous $D_m$ term.

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