Gauge-Mediated Supersymmetry Breaking: 
Introduction, Review and Update

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Recent progress in the gauge-mediated supersymmetry breaking is reviewed, with emphasis on the theoretical problems which gauge-mediated models so successfully solve, as well as those problems which are endemic to the models themselves and still beguile theorists today. (Talk given at the 5th International Conference on Supersymmetries in Physics (SUSY-97), Philadelphia, PA, May 27-31, 1997.)

For all the theoretical successes of supersymmetry (SUSY), its one overriding problem is, and has been, that there is no direct experimental evidence for it. But rather than throw SUSY out, we can try to retain some of its desirable properties (e.g., solution to the hierarchy/naturalness problems, unification of the gauge couplings) by building theories with *spontaneously broken* (or “softly-broken”) $N = 1$ SUSY. The central problem of SUSY would then appear to be “How is SUSY broken?” For many years this was the prime focus of research into SUSY. Two realizations have pushed that question somewhat to the sidelines. First, new exact results in SUSY have led to large numbers of new models in which SUSY is spontaneously broken. In the past three years alone, the number of such models has increased exponentially [1]. Now it appears that spontaneous SUSY-breaking is not a peculiar behavior exhibited by only a few special models, but is in fact rather generic. Second, it was realized many years ago [3] that if the mechanism for SUSY-breaking coupled too closely to the Standard Model (SM) spectrum, that spectrum would have to exhibit certain sum rules, namely $\text{STr} M^2 = 0$, which holds for sets of states with identical conserved quantum numbers. This relation implies that one of the $u$-type squarks must be lighter than the $u$-quark, in clear contradiction to experiment. However there is a way out of this phenomenological disaster, for this equation holds only at tree level, and only for renormalizable theories. The possibility then exists to break SUSY dynamically, but in some sector which only couples to the SM sector via loops or via non-renormalizable operators. We then refer to the SUSY-breaking sector as “hidden,” the SM sector as “visible,” and the intermediate states, which appear in loops or are integrated out to produce the non-renormalizable interactions, as “messengers” or “mediators.”

The banishment of SUSY-breaking to some mysterious “hidden sector” has an immediate phenomenological consequence. It is no longer the mechanism of spontaneous SUSY-breaking itself that prescribes the form which soft SUSY-breaking will take in the visible sector, but rather it is the mechanism for mediating SUSY-breaking that ultimately determines how SUSY will manifest itself in our colliders. The central problem of SUSY has changed. It is now “How does the Standard Model find out about SUSY-breaking?”

1. Communicating SUSY-breaking

There have evolved several approaches to dealing with the question of how SUSY-breaking is communicated to the SM. The first is simply to admit ignorance and to parametrize that ignorance with the most general effective Lagrangian consistent with the symmetries of the SM and with softly-broken $N = 1$ SUSY. A special case of this most general Lagrangian is the one which is minimal, that is the one which requires the least extension of the SM in order to accommodate the broken SUSY. This is the Minimal Supersymmetric Standard Model (MSSM). I call this approach...
“Don’t ask, don’t tell.” The “don’t ask” part is obvious, for we simply don’t worry about the exact nature of how SUSY-breaking is mediated. The “don’t tell” is the unfortunate part, for even in the MSSM (emphasis on “minimal”) there are 106 unknown parameters: 26 masses, 37 angles and 43 phases. With so many unknowns, extracting unequivocal experimental predictions for SUSY has proven to be a risky undertaking.

But already in this very general framework, we can begin to see that some organizing principle must be in place. With arbitrary, random parameters the MSSM has serious problems with flavor (and also CP) violation. Put simply, once the angles which rotate $U^{\dagger}U \neq 1$, leading to flavor-changing neutral currents (FCNCs).

These FCNCs show up in a variety of processes, though always at loop order: $K^0 - \bar{K}^0$ mixing, $\mu \rightarrow e\gamma$, etc. Consider the first of these. Requiring that the SUSY contributions to $\Delta m_K$ be smaller then the its experimental value yields the well-known contraint:

$$\text{Re} \left\{ \frac{\tilde{A}_L^2}{\tilde{m}_Q^2} \left( \frac{\delta m^2_{Q}}{\tilde{m}_Q^2} \right)^2 \right\} \lesssim 5 \times 10^{-9} \text{GeV}^{-2}$$

(1)

where $\tilde{m}$ is some average squark mass, $\delta m^2_Q$ is the $\tilde{d}_L - \tilde{\bar{d}}_L$ mass difference, and $\tilde{A}$ is a function of the angles which rotate $U \rightarrow \tilde{U}$. For generic $U$ and $\tilde{U}$ the average value for $\tilde{A}^2 \simeq 1/20$.

The constraints on CP violation in the kaon system are even stronger. Demanding that the SUSY contributions to $\epsilon_K$ not exceed its experimental value yields:

$$\delta \left| \frac{\tilde{A}_L^2}{\tilde{m}_Q^2} \left( \frac{\delta m^2_{Q}}{\tilde{m}_Q^2} \right) \left( \frac{\delta m^2_d}{\tilde{m}_Q^2} \right) \right| \lesssim 10^{-13} \text{GeV}^{-2}$$

(2)

where $\epsilon$ is the appropriate squark phase. Using Eq. (8) with $\delta \sim \mathcal{O}(1)$, $\tilde{A}^2 \simeq 1/20$ and $\tilde{m} = 500 \text{GeV}$, one finds that $m_{\tilde{g}} - m_{\tilde{d}} \lesssim 200 \text{MeV}$. Or for $\delta m^2 / \tilde{m}^2 \sim \mathcal{O}(1)$, we must have $\tilde{m} \gtrsim 700 \text{TeV}$! Thus FCNC and CP constraints demand either strong mass degeneracies among the scalars, or scalars so heavy that they decouple from the relevant processes altogether.

(There are at this time several competing proposals for solving this flavor/CP problems, including enforced mass degeneracies, decoupling of sparticles, and “alignment” in which $U = \tilde{U}$ and so $\mathcal{A} = 0$. I will only be discussing the first of these three possibilities in this talk.)

Within a most general Lagrangian, one has no way to approach the FCNC/CP problems; one has to consider specific mechanisms by which SUSY-breaking may be communicated from the hidden to the visible sector. The canonical method is supergravity. Local SUSY, a.k.a. supergravity, mixes the hidden and visible sectors through gravitational interactions (the mixing goes to zero as $M_{Pl} \rightarrow \infty$). The mixing terms, being non-renormalizable, suppress the scale of SUSY-breaking in the visible sector from that in the hidden sector. For example, for an O’Raifeartaigh-type breaking in the hidden sector at some scale $\sqrt{F}$, the apparent scale of SUSY-breaking in the visible sector will be $F/M_{Pl} \ll \sqrt{F}$.

The story of supergravity (SUGRA) is well-known. Mass universalities seem to naturally fall out without inputting them, solving both the FCNC and CP problems. One is left with only 5 free parameters at the SUGRA scale: a common scalar mass $m_0$, gaugino mass $M_{1/2}$, $A$-term $A_0$, $B$-term $B_0$, and of course $\mu$.

The truism that usually accompanies these universalities is that “gravity is flavor-blind.” I agree. Gravity has no reason to arrange its interactions so that they are diagonal in the same basis in which the Higgs couples to the fermions — and for that very reason it is clear that any basis chosen by gravity will differ from the Higgs interaction basis, leading to mass non-universalities and FCNCs. Even if gravity chooses no preferred ba-
sis at tree level, non-renormalizable terms in the Kähler potential, corrections coming from string interactions, renormalization group flow, and any of a dozen other possible effects will choose a preferred basis for the communication of SUSY-breaking, leading to disaster. In particular, if there is any source of flavor physics (e.g., any interactions which differentiate between $d$- and $s$-quarks) between the Planck scale and the weak scale, the mass universalities will be destroyed.

It would seem preferable for the mass degeneracies among the squarks and sleptons, rather than be accidents, to be guaranteed by the nature of the mediation mechanism. In particular, one would like the communication mechanism to respect global symmetries under which the various $u$-type quarks are identical, likewise for the various $d$-type quarks and the various leptons.

We don’t have to go far to find such symmetries, for they are the gauge symmetries of the SM. If the scalar soft masses were functions only of the gauge charges of the individual particles, universality would be automatic. And if the scale at which the communication of SUSY-breaking takes place is well below the Planck scale, then those Planckian “corrections” which upset universality in SUGRA cannot disrupt it here.

2. Gauge Mediation

The defining principle of gauge mediation, and the reason it solves the flavor problem, is that the SUSY partners of SM fermions receive the dominant part of their masses via gauge interactions.

The canonical model for gauge mediation was developed in the early 1980’s [7]. Suppose there exist some set of states (superfields) $\hat{\psi}$ with SM couplings but which are not part of the MSSM spectrum. Since the fermionic components of the $\hat{\psi}$ superfields must be heavy and not contribute to the SM anomalies, take them to be vectorlike with respect to the SM gauge interaction (i.e., $\bar{\hat{\psi}}\hat{\psi}$ is an SM singlet). In the part of the superpotential responsible for mediation, couple the $\hat{\psi}$ fields to a new singlet superfield $\hat{X}$:

$$W_M = \lambda \hat{X}\bar{\hat{\psi}}\hat{\psi} + \cdots$$ (3)

where $\hat{X}$ receives both $A$ and $F$ vevs. (Alternatively, the $A$ vev of $\hat{X}$ can be replaced with an explicit mass.) The source of the $F$ vev in the hidden sector is essentially irrelevant for most of the phenomenology and we take it as simply given; the exception is in the couplings of the gravitino to be discussed in Section 4.

$W_M$ of Eq. (3) induces masses for the fermionic components of $\hat{\psi}$ (which I denote $\hat{\psi}$) setting the overall scale for the messenger sector:

$$M \equiv m_\psi = \lambda X$$ (4)

while the scalar mass matrix has the form:

$$m^2_{\psi,\bar{\psi}} = \begin{pmatrix} \lambda^2 X^2 & \lambda F_X \\ \lambda F_X & \lambda^2 X^2 \end{pmatrix}.$$ (5)

The eigenvalues of Eq. (5) are:

$$m^2_{\psi,\bar{\psi}} = M^2 \pm \lambda F_X = M^2 (1 \pm x)$$ (6)

where $x \equiv \lambda F_X / M^2$; $x < 1$ in order to have positive squared masses. We define one more scale:

$$\Lambda \equiv \frac{\lambda F_X}{M}$$ (7)

which, as we will see, controls the weak scale.

Because the $\hat{\psi}$ fields are charged under the SM gauge groups, the gauginos of the MSSM can receive masses through loops of these new fields. In particular, the mass matrix of Eq. (5) contributes to the gaugino masses at 1-loop. Gaugino $\lambda_i$ of SM group $G_i$ receives a mass:

$$M_{\lambda_i}(M) = \frac{\alpha_i}{4\pi} g(x) T_i(\psi) \Lambda$$ (8)

where $T_i(\psi)$ is the Dynkin index of the representation of $\psi$ under $G_i$ and $g(x) = 1 + x^2 / 6 + \mathcal{O}(x^4)$.

The scalars of the MSSM do not receive soft masses until 2-loop order. For scalar $\phi$:

$$m^2_\phi(M) = 2 f(x) \Lambda^2 \sum_i \left( \frac{\alpha_i}{4\pi} \right)^2 C_i(\phi) T_i(\psi)$$ (9)

where $C_i(\phi)$ is the quadratic Casimir of the representation of $\phi$ under $G_i$ and $f(x) = 1 + x^2 / 36 + \mathcal{O}(x^4)$. Recall that there are also $D$-term contributions to scalar masses not included here. [An aside on conventions: I am using an $SU(5)$ normalization for $\alpha_1 = \frac{5}{3} \alpha_Y$ and so $T_1(\phi) = C_1(\phi) = \frac{5}{6} Y^2$ for any field $\phi$. For non-abelian $G_i = SU(n)$,}
\( T_i(\varphi) = \frac{1}{2} \) and \( C_i(\varphi) = \frac{x}{\Lambda} \) for \( \varphi \) in the fundamental representation.] For the oft-cited case in which \((\psi, \bar{\psi})\) are \(N\) pairs of \(\mathbf{5}, \mathbf{\bar{5}}\) of \(SU(5)\) and \(x \ll 1\), we notice that gaugino masses scale as \(N\), while scalar masses scale as \(\sqrt{N}\).

Finally, the trilinear soft terms (A-terms) arise at 2-loops. Since they have mass dimension-1, they are small compared to the rest of the soft masses and thus one can take \(A(M) \approx 0\). The case of the bilinear, dimension-2 \(B_{\mu}\) term will be discussed later.

One success of this particular mechanism for gauge mediation is that while gaugino masses arise at 1-loop, scalar mass-squareds arise at 2-loops. Thus gaugino and scalar masses are the same order in \(\alpha\). This is a noteworthy, because the simplest models of gauge mediation (those without the messenger fields) typically give masses to scalars at lower order than to gauginos, producing models with ultra-light gluinos.

One should make note of the scales that play a role in gauge mediation. Because LEP constrains the selectron mass \(m_{\tilde{e}} > 45\) GeV, then \(\Lambda > 30\) TeV. If by some fine-tuning argument we demand for gluinos that \(M_3 < 1\) TeV, then \(\Lambda \lesssim 120\) TeV. (These assume one pair of \(5 + \bar{5}\) messenger fields.) If all couplings in the problem are \(\mathcal{O}(1)\) then all mass scales will be \(\mathcal{O}(\Lambda)\), far below the Planck scale. Thus there will be no problems induced by supergravity corrections. (In more general models, the scales can differ greatly from one another. Then the requirement that supergravity corrections be small translates into the bounds \(F_X \ll m_Z M_{Pl}\) and \(M \lesssim 10^{15}\) GeV.)

### 3. Dine-Nelson Models

In the last few years, attention has been drawn back to gauge mediation as a viable alternative to supergravity mediation. Much of that renewed interest has been sparked by a series of models proposed by Dine, Nelson, and collaborators \[\text{[8]}\]. Because these models demonstrate both the successes and failings typical of models of gauge mediation, I will highlight their structure briefly.

Dine-Nelson models are divided into a tower of sectors, beginning with the hidden sector in which SUSY is broken spontaneously by some strong gauge dynamics. One often locates some non-anomalous global symmetry of the hidden sector which can be weakly gauged. That “messenger group,” typically a \(U(1)\), then communicates SUSY-breaking to a set of fields, \(\varphi^\pm\), which are SM singlets, giving them negative squared masses. The \(\varphi^\pm\) then couple to the messenger singlet \(X\) through terms in the superpotential:

\[
W_M = k \hat{X} \hat{\varphi}^+ \hat{\varphi}^- + \lambda \hat{X} \psi \bar{\psi} + k' \hat{X}^3. \tag{10}
\]

Setting \(\lambda = 0\) for now, the potential \(V(\varphi^\pm, X)\) is minimized when \(\langle \varphi^+ \varphi^- \rangle, \langle X \rangle\) and \(\langle F_X \rangle\) are all non-zero. Keeping that solution once \(\lambda\) is turned back on, the mass matrix of Eq. (3) is reproduced and SUSY-breaking is communicated to the MSSM through gauge interactions.

What are the successes of the gauge mediation approach? First, FCNCs are absent because of the scalar mass degeneracy. Second, scalar and gaugino masses are roughly the same size. Third, electroweak symmetry breaking occurs radiatively just as it does in supergravity models. And last, the models are highly predictive, with the entire spectrum of soft masses determined (to a good approximation) from just one input: \(\Lambda\).

### 4. Open Theoretical Issues

Despite their successes, the specific models of gauge mediation in the last section are open to a number of possible criticisms.

The global minimum of the scalar potential breaks color and not SUSY. The minimization performed in the last Section of the messenger potential \(V(\varphi^\pm, X)\) was not really correct for \(\lambda \neq 0\). With non-zero \(\lambda\), the minimum of the potential \(V(\varphi^\pm, X, \psi, \bar{\psi})\) occurs at \(\langle \psi \bar{\psi} \rangle = -\frac{4}{3} \langle \varphi^+ \varphi^- \rangle\) and \(\langle F_X \rangle = \langle X \rangle = 0\). That is, \(\psi\) does not find out about SUSY-breaking, but instead gets a non-zero vev, breaking \(SU(3) \times U(1)\). The minimum in which SUSY is broken is only local, not global \[\text{[9]}\].

There have been a number of suggestions for circumventing this problem. There are “cosmological” solutions to the problem, i.e., perhaps the present universe exists in a long-lived, metastable vacuum in which QCD and QED are preserved.
but SUSY is broken. There are also particle physics solutions. For example, we could add explicit masses for the ψ messenger fields to push their vevs to the origin \( \langle 10 \rangle \). Such masses may seem ad hoc but as long as they don’t have to sit at any one special scale, they are not unnatural. (In fact, they can sit comfortably at any scale between 100 TeV and \( 10^{15} \) GeV.) If the messenger group is a \( U(1) \), it is also possible to add extra matter which is chiral with respect to it \( \langle 11 \rangle \). Such extra matter can push the position of the minimum back to where we want it. However such matter can induce Fayet-Iliopoulos terms in the hidden sector, in which case it is known that there is no way to protect the messenger \( U(1) \) from mixing with hypercharge and destabilizing the visible sector \( \langle 11 \rangle \) (see below).

**Messenger \( U(1) \) interactions can be dangerous.**

For \( U(1) \) gauge groups, the field strength \( F_{\mu\nu} \) is gauge-invariant. Thus for theories with two (or more) \( U(1) \) groups, the gauge kinetic pieces of the Lagrangian can mix: \( \mathcal{L} \sim F_{\alpha\mu\nu} F_{\beta\mu\nu} \) for \( U(1)_a \times U(1)_b \). Because such terms are renormalizable, they can be induced by Planck-scale physics. Unless there is a symmetry to prevent such terms, one should in fact expect them to be present. And in SUSY, if the gauge kinetic pieces mix, then the \( D \)-terms must also.

Identify hypercharge as one of the \( U(1) \) factors and assume the other is in the hidden sector, as in the model of Section \( \langle 12 \rangle \). If (1) kinetic mixing occurs between the two \( U(1) \)'s, and (2) the \( D \)-vevs of the hidden sector \( U(1) \) are of order the scale of SUSY-breaking in that sector, then the large hidden sector \( D \)-term will be communicated to the scalars of the MSSM, pulling their masses up to the scale of hidden sector SUSY-breaking \( \langle 12 \rangle \).

There are two ways to get around this problem. The first is to find a charge-conjugation symmetry which acts on only one of the \( U(1) \)'s, namely \( C : A_\mu^H \rightarrow -A_\mu^H \) while \( C : A_\mu^H \rightarrow A_\mu^H \). Such a symmetry will forbid kinetic mixing, particularly if it is a gauged discrete symmetry. The second possibility is to work only with non-abelian messenger groups for which kinetic mixing cannot occur.

The \( \mu \)-problem is worse than usual. The \( \mu \)-problem of the MSSM is familiar. If \( W = \mu H_U H_D \) is a SUSY- and \( G_{SM} \)-invariant mass term, why is \( \mu \sim m_Z \) instead of \( \mu \sim M_{Pl} \)? We know that \( \mu \sim m_Z \) because it contributes to the Higgs potential and thus to the \( Z^0 \) mass:

\[
\mu^2 = \frac{m^2_{H_D} - m^2_{H_U} \tan^2 \beta}{\tan^2 \beta - 1} - \frac{1}{2} m_Z^2 \tag{11}
\]

where all the masses on the RHS are \( \mathcal{O}(m_Z) \). Within the context of supergravity, the most promising solution is the Giudice-Masiero mechanism \( \langle 12 \rangle \) in which the hidden and visible sectors mix through a non-minimal Kähler potential, \( K = K_0 + \frac{1}{M_{Pl}} X^\dagger H_U H_D + \text{h.c.} \), where \( K_0 \) is the canonical piece and \( X \) is some hidden sector field with \( F_X \sim m_Z M_{Pl} \). Then:

\[
\int d^4 \theta \frac{X^\dagger}{M_{Pl}} \dot{H}_U \dot{H}_D = \int d^2 \theta \frac{F_X}{M_{Pl}} \dot{H}_U \dot{H}_D, \tag{12}
\]

which contributes to the superpotential with coefficient \( F_X/M_{Pl} \sim m_Z \) and is the usual \( \mu \)-term. In gauge-mediated models this mechanism cannot work for the same reason that other supergravity contributions are small: \( F_X/M_{Pl} \ll m_Z \).

Gauge-mediated models also have a second, related problem which does not show up in supergravity. The dimension-2, bilinear soft term \( B_\mu H_U H_D \) does not arise until 3-loops or beyond. It is therefore, like the \( A \)-term, essentially zero. This is not phenomenologically feasible since the pseudoscalar Higgs of the MSSM gets a mass \( m_A^2 = 2 B_\mu / \sin 2 \beta \). As \( B_\mu \rightarrow 0 \), \( m_A \rightarrow 0 \) and \( A^0 \) becomes an axion.

The oldest solution to the \( \mu \)-problem is to extend the MSSM to the NMSSM, which is the MSSM plus a singlet with superpotential

\[
W = \lambda_H \hat{S} \hat{H}_U \hat{H}_D + \kappa \hat{S}^3 \tag{13}
\]

Then \( \langle S \rangle \) provides the \( \mu \)-term. An obvious candidate for that singlet is the field \( X \) which appears in \( W_M \langle 14 \rangle \). Here \( \langle X \rangle \) would provide a \( \mu \)-term and \( \langle F_X \rangle \) would provide a \( B_\mu \) term, seemingly as desired. But since \( \langle X \rangle \gg m_Z \), a solution to the \( \mu \)-problem clearly requires \( \lambda_H \ll 1 \). Then

\[
\mu = \lambda_H X \tag{14}
\]

\[
B_\mu = \lambda_H X^2 \gg \lambda_H^2 X^2 \tag{15}
\]
Thus $B_\mu \gg \mu^2$. If we choose $\mu \sim m_Z$ then $B_\mu$ is huge, causing the Higgs potential to become unbounded from below. If we choose $B_\mu \sim m_Z$ then $\mu$ is tiny, leading to light charginos that would have been found at LEP.

One could try to avoid the problems intrinsic to $X$ by introducing a new singlet $S$ just to solve the $\mu$-problem [13]. Unfortunately, being a gauge singlet, $S$ does not receive a very large soft mass and so its physical component is very light. Another way to see this is to notice that the superpotential of Eq. (1) has an $R$-symmetry which is only broken at 2-loops by small $A$-terms. Thus the light field is really an $R$-axion.

Finally, it has been suggested that non-renormalizable interactions in the superpotential could conspire to produce $\mu \sim m_Z$ [3]. Such models again yield $B_\mu \ll m_Z^2$.

There is still a CP/electric dipole problem. Despite the fact that the CP problem in the kaon system has been resolved by the mass degeneracies, there remains another CP problem which persists. In the simplest models of gauge mediation, there are 4 non-zero CP-violating phases beyond those of the SM: arg($\mu$), arg($A$), arg($B_\mu$) and arg($M_3$). Of these only two combinations are physical [4]:

$$\Phi_A = \text{arg}(A^* M_3) \quad \text{and} \quad \Phi_B = \text{arg}(B_\mu^* \mu M_3). \quad (16)$$

These phases contribute to electric dipole moments of quarks and leptons, and in turn, nuclei. One finds for the neutron electric dipole moment:

$$d_N \approx 2 \left( \frac{100 \text{ GeV}}{m} \right)^2 \sin \Phi_{A,B} \times 10^{-23} \text{ e cm} \quad (17)$$

for some generic squark mass $m$. Experimentally $d_N < 1.1 \times 10^{-25} \text{ e cm}$. Thus $O(1)$ phases are only allowed if $m \gtrsim 1 \text{ TeV}$. Any solution which doesn’t include very heavy squarks must compensate by finding some way to enforce small phases.

5. The Minimal Messenger Model

Recall now the earlier discussion of the $\mu$- and $B_\mu$ problems. It should be clear that in gauge-mediated models, the solutions to these two problems are intimately connected. And unfortunately we usually can solve one only at the expense of the other. Now we will show that certain types of gauge-mediated models actually solve both problems simultaneously, while at the same time resolving the CP/electric dipole problem and being highly predictive.

Start from a gauge-mediated model in which the $\mu$-problem has been solved at the expense of small $B_\mu \ll \mu^2$. As was said before, finite (threshold) contributions to $B_\mu$ only come in at 3-loops and thus are very small. However the divergent contributions receive log enhancement. In supergravity, such an enhancement is huge; here it is relatively small, but enough to lift the axion to experimentally allowed masses [13].

There are two symmetries of the MSSM which try to enforce $B_\mu = 0$: a Peccei-Quinn symmetry, broken by non-zero $\mu$, and an $R$-symmetry, broken by gaugino masses. Thus we expect contributions to $B_\mu$ proportional to $\mu M_\lambda$. This in fact happens through the 1-loop RGE’s for $B_\mu$ which have an approximate solution (for $M = 100 \text{ TeV}$):

$$B_\mu(m_Z) \approx \mu M_\lambda(M) \left[ -0.12 + 0.17 y_t^2 \right] \quad (18)$$

$$\sim \alpha_2 \mu M_2. \quad (19)$$

For $\mu \sim M_2 \sim m_Z$ we can write $\tan \beta \sim m_Z^2 / B_\mu \sim \alpha_2^{-1} \sim 30$. The “axion” mass is $M_A \approx 400 \text{ GeV}$. (More precise calculations including the full effective potential and threshold effects yields [10] $\tan \beta \approx 50$.) This highly predictive version of gauge-mediation, in which $B_\mu$ is set to zero at the messenger scale and is therefore no longer a free parameter, has been called the Minimal Messenger Model (MMM) [15,16].

Such large values for $\tan \beta$ will strike those familiar with supergravity as unnatural, following the arguments of Ref. [17]. However those arguments work only for models with one mass scale. Such is not the case here, where the mass scales of interest are related but very different: $B_\mu \sim \alpha_2 \mu M_2$. The hierarchy in the vevs is then simply a reflection in the hierarchy of the mass scales.

The MMM has one other great advantage. Because $B_\mu(M) = A(M) = 0$, then too $\Phi_A(M) = \Phi_B(M) = 0$. While the first relation is modified by RGE’s when running down to the weak scale,
the second is not. This is because:
\[
\frac{dB^\mu}{d\log Q} \propto M_i^5 \quad \text{and} \quad \frac{dA^*}{d\log Q} \propto M_i^7
\]
which renders \( \Phi_{A,B} = 0 \) a stable fixed-point. Thus the CP/electric dipole problem has been solved by finding a natural way in which to suppress the new SUSY phases without having to enforce very heavy particle spectra.

6. Direct Transmission Models

In trying to overcome some of the difficulties of the Dine-Nelson models highlighted in Section 3, some natural questions arise, such as “Can we live without a messenger interaction?” and “Can \( \psi \) live in the SUSY-breaking sector itself?” The answer to both questions is “yes” provided the SUSY-breaking sector possesses an anomaly-free global symmetry large enough to embed the entire SM gauge structure and which is not broken concurrent to SUSY breaking. Such models are called “direct transmission” models.

Finding such large global symmetries in the hidden sector seems to be quite a difficulty, for it generally requires very large dynamical SUSY-breaking groups such as \( SU(15) \) or \( SU(7) \times SU(6) \). The large size of the dynamical groups feeds back into the visible sector because we want the \( \psi \) to transform under the fundamental representation of the dynamical groups. Thus a model whose dynamical group is \( SU(n) \) containing a global \( SU(5) \) will have \( n \) copies of \( 5 + \overline{5} \). If \( \sqrt{F} \sim M \sim 50 \text{ TeV} \), then perturbativity up to the Planck scale requires \( n < 4 \), which is too small.

The solution to this conundrum is obvious: move the messenger mass scale far abovet he weak scale, and consequently, far above the SUSY-breaking scale \( \sqrt{F} \). The means to get this is less obvious [13]. In most Dine-Nelson type models, the scalar potential of the messenger sector has no flat directions, and so on minimization one finds \( F_X \sim X^2 \) for any field \( X \) participating in SUSY breaking. However, along a flat direction the potential doesn’t turn up until very large values of \( X \) so that at the minimum \( F_X \ll X^2 \). This precisely the desired behavior. One caveat however: we need to ensure that the supergravity contributions remain small compared to the gauge-mediated contributions. Thus we must have:
\[
\langle X \rangle \lesssim 10^{15} \text{ GeV}
\]
\[
\frac{F_X}{M_{Pl}} \lesssim \sqrt{n} \frac{\alpha}{4\pi} \frac{F_X}{X}
\]
for dynamical \( SU(n) \) with gauge coupling \( \alpha \).

Direct transmission models have most, if not all, of the following properties: (i) Gauge coupling unification occurs just as in the MSSM. In particular, there is no loss of perturbativity before the unification scale. (ii) There are no gauge singlets in the model. (iii) The SUSY-breaking, QCD/QED-preserving minimum is global. (iv) The squarks are no longer quite so heavy compared to the sleptons. The large amount of running to go from \( Q = M \) to \( m_Z \) washes out the largest mass hierarchies among the sparticles. (v) The gravitino has a mass \( m_{3/2} \sim \text{few GeV} \). This may be a serious cosmological problem (see Section 8). (vi) The detector signatures will very closely resemble those of supergravity models. In particular, the NLSP decay to gravitinos occurs outside the detector (see Section 8). Finally, (vii) \( \text{STr} M^2 > 0 \) where the supertrace is only over the set of \( \psi + \overline{\psi} \) messenger fields.

The impact of that last statement is only appreciated when one considers the 2-loop running of the MSSM scalar masses. There one finds a correction to the soft scalar masses of the form [18]:
\[
\delta m_\phi^2 = -\frac{\alpha^2}{4\pi^2} C_i(\phi) T_i(\psi) \text{STr} M^2 \log \frac{\Lambda_{UV}}{m_\psi}.
\]
In the direct transmission models, the messenger sector contains a number of \( \psi \)-fields and there may be a sizeable hierarchy which develops among them. Then \( \Lambda_{UV} \) is to be interpreted at the largest mass in the messenger sector. If \( \text{STr} M^2 > 0 \) and some \( m_\psi \ll \Lambda_{UV} \) then \( \delta m_\phi^2 \) will push MSSM scalars to negative squared masses. Such is the case in the models of Refs. [18], but not that of Ref. [19]. For all three models, the bulk of the messenger sector sits near \( 10^{15} \text{ GeV} \) while a few fields sit near \( 100 \text{ TeV} \); in the last case, however, those light fields are eaten as the flavor symmetry breaks down from \( SU(10) \) to \( SU(5) \).
7. Phenomenology and the Gravitino

The role of the gravitino in SUSY models is well-known. In global SUSY, spontaneous breaking produces a massless (spin-1/2) goldstino with derivate couplings to the SUSY current. This is expressed in the SUSY generalization of the Goldberger-Treiman formula [20]:

\[ \mathcal{L} \sim \frac{1}{F} \left( \bar{X} \gamma^\mu \sigma^{\mu\nu} \partial_\nu \tilde{G} F^A_{\mu\nu} + \psi_L \gamma^\mu \gamma^\nu \partial_\mu \tilde{G} \partial_\nu \phi \right). \]

Here \( F \) is the \( F \)-term responsible for the original SUSY-breaking in the hidden sector; it can be thought of as being the largest \( F \)-term in the theory. The decay width of sparticles into their partners plus gravitinos can then be calculated. For example, the width for a bino, \( B \), decaying into a photon and gravitino is given by:

\[ \Gamma(B \to \gamma \tilde{G}) = \frac{\cos^2 \theta_W}{8\pi} \frac{m_B^2}{F^2} \]

which becomes larger as \( F \) becomes smaller.

Once SUSY is elevated to a local symmetry (i.e., supergravity), the goldstino is eaten via the super-Higgs mechanism by the massless (spin-3/2) gravitino. The gravitino acquires longitudinal components and a mass: \( m_{3/2} \approx F/M_{Pl} \). Since gravity is so weak, it is only the longitudinal components of \( \tilde{G} \) that interact. Therefore the results which were derived for the goldstino hold equally well for the gravitino.

In supergravity models one rarely worries about \( \tilde{G} \). It is too weak to be produced directly in experiments, and there is no reason why it should be lighter than any other SUSY partner so that other states decay into it. (For my purposes here, I am always assuming that there is a discrete symmetry, like \( R \)-parity, which makes the lightest SUSY particle (LSP) absolutely stable.) In gauge-mediated models the gravitino is light, roughly \( 1 \text{ eV} \) to \( 1 \text{ GeV} \), making it the LSP. It is still too weak to be directly produced in experiments, but as the LSP, all other SUSY particles must eventually decay into it. The phenomenology of gauge-mediation, which is otherwise so much like that of supergravity, has a new component, the search for decays into gravitinos.

There are two central questions which arise in studying gravitino phenomenology [21][22]. First, what is the NLSP (the next-to-lightest SUSY particle)? Even with light gravitinos, the SUSY states will dominantly decay via their strong and electroweak interactions, until the only sparticles left are the NLSPs. The NLSP, having no other route for its decay, eventually goes to the gravitino plus some other particle(s) whose identity relies heavily on the type of NLSP present. The second question is, what is the decay width (or decay length) of the NLSP into gravitinos? Since the NLSP decay length scales as \( F^2 \), a measurement of that length is a direct measurement of SUSY-breaking in the hidden sector!

The answer to the first question may be model-dependent, but is usually one of only a few possibilities. For models with only one pair of messenger fields, and for tan \( \beta \) small to intermediate, the NLSP is a neutralino, \( N_1 \), which is itself usually bino-like. For larger multiplicities of messenger fields, the NLSP(S) are the RH sleptons. But as tan \( \beta \) increases, the \( \tilde{\tau} \) becomes the sole NLSP. For bino NLSP, we can expect decays most often to \( \tilde{G} + \gamma \); for slepton NLSP we can expect decays to the partner lepton, and in particular, \( \tau \)-leptons.

The answer to the second question dictates whether or not the NLSP decays inside or outside the detector, and if inside, whether the decay length is long enough to be reconstructed. The various possibilities are given in the following table adapted from Ref. [22].

| NLSP          | Decay Length     | Signal          |
|---------------|------------------|-----------------|
| \( N_1 \approx \tilde{B} \) | Prompt           | \( \gamma \gamma + \not{E}_T \) |
| \( \tilde{\ell} = \tilde{e}, \tilde{\mu}, \tilde{\tau} \) | 2nd Vertex       | \( \gamma \gamma + \not{E}_T \) |
|               | Outside          | \( \not{E}_T \) |
|               | Prompt           | \( \ell \ell + \not{E}_T \text{ w/ kinks} \) |
|               | Outside          | Heavy Leptons   |

Here “prompt” decays are too close to the vertex to differentiate, \( ^{2\text{nd}} \text{vertex} \) refers to differentiated second vertices at which the decay to gravitinos occurred, and “outside” means that the decay took place outside of the detector. The signals are self-explanatory apart from the following note: “kinks” are sudden turns in the track of a charged particle, in this case occurring...
where the invisible $G$ is being emitted. “Heavy leptons” means that the track will be charged, but not jetty, and will reconstruct to have a mass far above normal lepton masses. Of course, it is also possible that more than one option in the table is realized. Either the decay length could put it on one of the boundaries in the table, or the lightest slepton and neutralino could be so close in mass that they prefer to decay to gravitinos rather than to one another.

8. Gauge Mediation and Cosmology

In any gauge-mediated model, the gravitino will be the LSP. Its mass, however, is model-dependent. In the Dine-Nelson models we can estimate the gravitino mass:

$$m_{3/2} \sim \frac{1}{M_{Pl}} \left( \frac{16\pi^2}{g_M} \Lambda \right)^{2} \sim 20 \text{ keV}$$

(25)

where I have assumed the messenger group coupling, $g_M$, in the last equality to be $O(1)$. For direct transmission models, the mass is larger, about 1 GeV. Finally, though there is no full model at present that does such, models in which the NLSP decays to gravitinos inside the detector will have $m_{3/2} < 1$ keV.

Using the standard techniques, the relic abundance of gravitinos present in the universe today can be calculated. For gravitinos whose abundances are not diluted by some mechanism, one needs $m_{3/2} < 2\hbar^2 \text{ keV}$ in order to avoid overclosing the universe. ($h$ is the Hubble constant in units of 100 km/sec/Mpc.) Thus the gravitinos which would be implied by observing NLSP decay would be cosmologically acceptable and may even be a useful source of dark matter.

In the mass range $1 \text{ keV} \lesssim m_{3/2} \lesssim 100 \text{ keV}$, such as in the Dine-Nelson models, an overabundance of gravitinos results from late NLSP decay. To wash out this overabundance one would like a period of late inflation, with the constraint that the reheating temperature, $T_R$, is less than $m_\sim \sim m_\lambda$ in order to avoid producing more of the NLSP which would again decay to gravitinos.

In the mass range $100 \text{ keV} \lesssim m_{3/2} \lesssim 5 \text{ GeV}$, such as one finds in direct transmission models, the overabundance of gravitinos is produced through scattering processes: $A + B \rightarrow C + \tilde{G}$. Once again a period of late inflation can wash out the excess gravitino abundances, but only as long as $T_R$ does not get so large as to reproduce the conditions of the gravitino production. In this case, that means $T_R < 10^8 m_{3/2}$. Finally for gravitinos with $m_{3/2} \gtrsim 5 \text{ GeV}$, the decay width $\Gamma(NLSP \rightarrow G)$ is very narrow. The NLSP doesn’t decay until after nucleosynthesis, destroying light nuclei by photofission. It seems very difficult to avoid the problems associated with a gravitino in this mass range.

There is one other cosmological concern one might have in gauge-mediated models which only arises when coupling the model to string theory. In string theory, there are fields called string moduli whose vevs parametrize the size of the compactified extra dimensions but whose potentials are flat before SUSY is broken. After SUSY breaking, the moduli get masses $m \sim m_{3/2}$.

In the high density, high temperature early universe the potential for the moduli gets additional contributions proportional to the Hubble constant, $H$, and the temperature, $T$, since both break SUSY. But the minimum of the moduli potential at finite $H$ and $T$ need not be the same as the minimum when $H = T = 0$. Thus when $H, T$ fall below $m_{3/2}$ the moduli begin falling towards their true minimum. However, since their couplings to matter are Planck-suppressed, there is little damping and the moduli begin oscillating about their minima with amplitudes of $O(M_{Pl})$. During the period of their oscillations, they contribute to the energy density of the universe through their $(\nabla \phi)^2$ kinetic energy.

If the potential is too shallow and the damping too small, the moduli are still oscillating today. That energy density would easily overclose the universe unless some means was found to inflate away the moduli. Such an inflation seems difficult to arrange for the moduli of the Dine-Nelson models, but may be possible for the heavier direct transmission moduli.

9. Conclusions

The study of gauge-mediated SUSY-breaking has flowered dramatically over the last year.
Though much of the early attention on these models focused on the single $ee\gamma + E_T$ event at CDF, the theoretical interest and justification for these models go well beyond any single experimental anomaly. Pragmatically, these models provide an alternative measure for testing experimental sensitivities to SUSY in all its guises. And they are simple to work with, with only 3 free parameters: $\Lambda$, $\tan\beta$ and $\log M$. One particularly attractive version of gauge mediation (the MMM) has only two, $\tan\beta$ being an output.

Ideally, gauge-mediated models provide an opportunity to do something which is usually very difficult, and that is to probe the physics of the hidden sector directly. And by solving the SUSY flavor problem, they allow interesting flavor physics to occur at scales not far above the weak scale without inducing large FCNCs.

Progress in this area seems to require effort in two directions right now: careful study of the phenomenology of these models at current and future experiment, with special attention paid to going beyond the minimal models; and continued attempts to build realistic models of gauge-mediation which work around the problems I have outlined here. If progress continues at the same rate this coming year as it has this past year, a review talk on gauge-mediation at SUSY-98 would be both very difficult and very exciting to give.

REFERENCES
1. See the talk of A. Nelson in this volume.
2. S. Dimopoulos and H. Georgi, Nucl. Phys. B193, 150 (1981).
3. L. Girardello and M. Grisaru, Nucl. Phys. B194, 65 (1982). L. Hall and L. Randall, Phys. Rev. Lett. 65, 2939 (1990).
4. S. Dimopoulos and D. Sutter, Nucl. Phys. B452, 496 (1995). See also the talk of H. Haber in this volume.
5. You know who you are.
6. M. Dine, A. Kagan and S. Samuel, Phys. Lett. B243, 250 (1990); F. Gabbiani, et al, Nucl. Phys. B477, 321 (1996).
7. M. Dine, W. Fischler and M. Srednicki, Nucl. Phys. B189, 575 (1981); S. Dimopoulos and S. Raby, Nucl. Phys. B192, 353 (1981); C. Nappi and B. Ovrut, Phys. Lett. 113B, 175 (1982); L. Alvarez-Gaumé, M. Claudson and M. Wise, Nucl. Phys. B207, 96 (1982).
8. M. Dine and A. Nelson, Phys. Rev. D48, 1277 (1993); M. Dine, A. Nelson and Y. Shirman, Phys. Rev. D51, 1362 (1995); M. Dine, et al, Phys. Rev. D53, 2658 (1996).
9. I. Dasgupta, B. Dobrescu and L. Randall, Nucl. Phys. B483, 95 (1997).
10. L. Randall, Nucl. Phys. B495, 37 (1997).
11. K. Dienes, C. Kolda and J. March-Russell, Nucl. Phys. B492, 104 (1997).
12. G. Giudice and A. Masiero, Phys. Lett. B206, 480 (1988).
13. G. Dvali, G. Giudice and A. Pomarol, Nucl. Phys. B478, 31 (1996).
14. For a nice review, see Y. Grossman, Y. Nir and R. Rattazzi, hep-ph/9701231.
15. K.S. Babu, C. Kolda and F. Wilczek, Phys. Rev. Lett. 77, 3070 (1996); M. Dine, Y. Nir and Y. Shirman, Phys. Rev. D55, 1501 (1997).
16. J. Bagger, et al, Phys. Rev. D55, 3188 (1997); R. Rattazzi and U. Sarid, hep-ph/9612464; F. Borzumati, hep-ph/9702307.
17. A. Nelson and L. Randall, Phys. Lett. B316, 516 (1993).
18. E. Poppitz and S. Trivedi, Phys. Rev. D55, 5508 (1997), Phys. Lett. B401, 38 (1997); N. Arkani-Hamed, H. Murayama and J. March-Russell, hep-ph/9701286.
19. H. Murayama, Phys. Rev. Lett. 79, 18 (1997).
20. P. Fayet, Phys. Rep. 105, 21 (1984).
21. S. Dimopoulos, S. Thomas and J. Wells, Phys. Rev. D54, 3283 (1996), Nucl. Phys. B488, 39 (1997).
22. S. Ambrosanio, et al, Phys. Rev. D54, 5395 (1996).
23. H. Pagels and J. Primack, Phys. Rev. Lett. 48, 223 (1982).
24. T. Moroi, H. Murayama and M. Yamaguchi, Phys. Lett. B303, 289 (1993).
25. A. de Gouvea, T. Moroi and H. Murayama, Phys. Rev. D56, 1281 (1997).
26. S. Park, in “Proceedings of the 10th Topical Workshop on Proton-Antiproton Collider Physics,” ed. by R. Raja and J. Yoh, AIP Press (1995).