MECHANISMS OF SUPERSYMMETRY BREAKING IN THE MSSM

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Preliminary Remarks. Gauge mediated supersymmetry breaking. Gravity mediated supersymmetry breaking. Anomaly mediated supersymmetry breaking. Gaugino mediated supersymmetry breaking. Braneworld supersymmetry breaking. Conclusions.

• Preliminary Remarks

This will be a somewhat theoretical review of models and mechanisms for generating soft explicit supersymmetry breaking terms in the MSSM. There won’t be much signal phenomenology except in a few illustrative cases. Also, I shall be somewhat antihistorical in first talking about gauge mediation and then coming to gravity mediation since my subsequent topics, i.e. AMSB, gaugino mediation as well as braneworld scenarios, connect more naturally with the latter.

Our Lagrangian can be decomposed \([1]\) as

\[
\mathcal{L}_{\text{MSSM}} = \mathcal{L}_{\text{MSSM}} + \mathcal{L}_{\text{SOFT}},
\]

\[
-\mathcal{L}_{\text{SOFT}} = \frac{1}{2}(M_1 \tilde{\lambda}_0 \tilde{\lambda}_0 + M_2 \tilde{\lambda} \cdot \tilde{\lambda} + M_3 \tilde{g}^a \tilde{g}^a + \text{h.c.}) + V_{\text{SOFT}}^{\text{SCALAR}},
\]
\[
V_{\text{SCALAR}}^{\text{SOFT}} = \sum_j \tilde{f}_j^*(\mathcal{M}^2_f)_{ij} \tilde{f}_i + (m_1^2 + \mu^2)|h_1|^2 + (m_2^2 + \mu^2)|h_2|^2 \\
+ (B \mu h_1 \cdot h_2 + h.c.) + \text{trilinear A terms.}
\] (1c)

The sfermion summation in (1c) covers all left and right chiral sleptons and squarks. The other scalars, namely the Higgs doublets \( h_{1,2} \), occur explicitly in the RHS. A direct observable consequence of (1) is the upper bound \([1]\) on the lightest Higgs mass

\[ m_h < 132 \text{ GeV}, \]

which is a ‘killing’ prediction of the MSSM.

Though \( \mathcal{L}_{\text{SOFT}} \) provides a consistent and adequate phenomenological description of the MSSM, it is ad hoc and ugly. One would like a more dynamical understanding of its origin. Supersymmetry has to be broken and spontaneous breakdown would be an elegant option. Unfortunately, if this is attempted with purely MSSM fields, disaster strikes in the form of the Dimopoulos-Georgi sumrule \([1]\):

\[ STr M^2_{\ell_i} + STr M^2_{\nu_i} = 0 = STr M^2_{u_i} + STr M^2_{d_i}, \] (2)

where \( STr M^2_j \equiv m_{\tilde{f}_j}^2 + m_{\tilde{f}_j}^2 - 2m^2_f \) in terms of physical masses and \( i \) is a generation index. Evidently, (2) is absurd since, for each generation, some sparticles are predicted to be lighter than the corresponding particles in contradiction with observation.

The way out of this conundrum is to postulate a hidden world of superfields \( \Sigma \) which are singlets under SM gauge transformations. Let spontaneous supersymmetry breaking (SSB) take place at a scale \( \Lambda_S \) in this hidden sector and be communicated to the observable world of superfields \( Z \) by a set of messenger superfields \( \Phi \) (Fig. 1) – characterized by some messenger scale \( M_m \). The induced soft supersymmetry breaking parameters in the observable sector get characterized by the particle-sparticle mass splitting \( \sim M_s = \Lambda_S M_m^{-1} \). The messengers could all be at the Planck scale (i.e. \( M_m = M_{PL} \)), but such need not be the case. They

\[ \text{Figure 1: The transmission of supersymmetry breaking.} \]
are, in fact, two broad categories of messenger mechanisms: (1) gauge mediation and (2) gravity mediation. In (1) the messengers are intermediate mass ($\geq 100$ TeV) fields with SM gauge interactions. In (2) they are near Planck scale supergravity fields inducing higher dimensional supersymmetry breaking operators suppressed by powers of $M_{PL}^{-1}$.

- **Gauge mediated supersymmetry breaking** [2, 3, 4]

  The messenger superfields here have all the MSSM gauge interactions. MSSM superfields, with identical gauge interactions but different flavours, are treated identically by the messengers; thus there are no FCNC amplitudes. Loop diagrams induce the explicit soft supersymmetry breaking terms in the MSSM. Loop diagrams, generating gaugino and scalar masses, are shown in Figs. 2a and 2b with $\{\phi, \chi\}$ and $\{Z, \psi\}$ being components of $\Phi$ and $Z$ respectively. Let $S$ be a generic hidden sector chiral superfield and $\{\Phi_i, \bar{\Phi}_i\}$ a set messenger chiral superfields\footnote{$\Phi_i$ and $\bar{\Phi}_i$ together form a vectorlike representation of $SU(5)$}, interacting via couplings $\lambda_i$ in the superpotential

  \[ W_{\text{mess}} = \sum_i \lambda_i S \Phi_i \bar{\Phi}_i. \]  

  SSB in the hidden sector is characterized by the auxiliary component VEV $\langle F_S \rangle$. A typical messenger mass is given by $M_m \sim |\lambda_i \langle S \rangle|$. Define

  \[ x_i \equiv \left| \frac{\langle F_S \rangle}{\lambda_i \langle S \rangle} \right|, \quad \Lambda \equiv \left| \frac{\langle F_S \rangle}{\langle S \rangle} \right|, \]  

  i.e. $M_m = \Lambda / x_i$. One can then show from the required positivity of the lowest eigenvalue of the messenger scalar mass matrix that $0 < x_i < 1$.

  \[ M_\alpha = \left( \frac{g_2^2}{16\pi^2} \right) \Lambda \sum_\alpha 2T_\alpha(R_i)g(x_i), \]  

  \[ m^2_{f, h} = 2\Lambda^2 \sum_\alpha (g_2^2/16\pi^2)^2 C_\alpha \sum_i 2T_\alpha(R_i)f(x_i). \]
Here \( T^a(\phi_i)T^b(\phi_i) = T^a(R_i)\delta^{ab} \) where the trace is over the representation \( R_i \) of \( \phi_i \) in the gauge group factor \( G_\alpha \) and \( C_\alpha \) is the quadratic Casimir \( \left( \sum_a T^a T^a \right)_{G_\alpha} \) of the latter. Moreover,

\[
\begin{align*}
\text{(6a)} & \quad g(x) = x^{-2}[(1 + x) \ln(1 + x) - (1 - x) \ln(1 - x)], \\
\text{(6b)} & \quad f(x) = x^{-2}(1 + x) \left[ \ln(1 + x) - 2Li_2 \left( \frac{x}{1 + x} \right) + \frac{1}{2}Li_2 \left( \frac{2x}{1 + x} \right) \right] + (x \leftrightarrow -x),
\end{align*}
\]

\( Li_2 \) being the dilogarithm. The behaviour of \( g(x) \) and \( f(x) \) in the region \( 0 \leq x \leq 1 \) is shown in Fig. 3. They are practically unity for a large range of \( x \). In this situation \( \sum_\alpha 2T^a(R_i) \) factorizes and becomes \( n_5^5 \) for \( SU(3)_C \) or \( SU(2)_L \) but \( \sum_i (Y_i/2)^2 = \frac{n_5}{3}n_5 \) for \( U(1)_Y \), where \( n_5 \) is the number of complete \( 5 \oplus 5 \) messenger representations of \( SU(5) \). Now one can write

\[
\begin{align*}
M_\alpha & \simeq \left( g_\alpha^2/16\pi^2 \right)n_5\Lambda \\
\text{(7a)} & \quad m^2_{\tilde{f}, M}(M_m) \simeq 2n_5^{-1} \left[ C_3M_3^2(M_m) + C_2M_2^2(M_m) + \frac{3}{5} \left( \frac{Y}{2} \right)^2 M_1^2(M_m) \right], \quad \text{(7b)}
\end{align*}
\]

where \( C_3 = \frac{4}{3} \) (0) for an \( SU(3)_C \) triplet (singlet) and \( C_2 = \frac{2}{3} \) (0) for an \( SU(2)_L \) doublet (singlet). To one loop the gaugino masses \( \text{(7a)} \) vary with RG evolution in the same way as \( g_\alpha^2 \), while the scalar masses \( \text{(7b)} \) are specified at an energy scale \( M_m \) corresponding to messenger masses. The trilinear coupling \( A \) parameters get induced at the two loop level and can be taken to vanish at the scale \( M_m \) — becoming nonzero at lower energies via RG evolution. The parameters \( \mu, B \) are kept free to implement the radiative EW breakdown mechanism, the validity of which implies the bounds

\[
50 \text{ TeV} < M_m < \sqrt{n_5} \times 10^{14} \text{ GeV}. \quad \text{(8)}
\]

The minimal GMSB model, called mGMSB, is characterized by the parameter set

\[
\{ p \} = \{ \Lambda, M_m, \tan \beta, n_5, \text{sgn} \mu \}. \quad \text{(9)}
\]

Linear RG interpolation of sfermion squarel masses from the boundary values of \( \text{(7b)} \) at the scale \( M_m \) to lower energies \( \sim \Lambda \) yield, with \( t_M = \ln M_m/\Lambda \), the one loop expressions.

Figure 3: The behaviour of (a) \( g(x) \) and (b) \( f(x) \)
\[ m^2_{\tilde{e}_R}(100 \text{ GeV}) = M_1^2(100 \text{ GeV}) \left[ 1.54n_5^{-1} + 0.05 + (0.072n_5^{-1} + 0.01)t_M \right] \]

\[ + s^2_W D, \quad (10a) \]

\[ m^2_{\tilde{e}_L}(100 \text{ GeV}) = M_2^2(100 \text{ GeV}) \left[ 1.71n_5^{-1} + 0.11 + (0.023n_5^{-1} + 0.02)t_M \right] \]

\[ + (0.5 - s^2_W)D, \quad (10b) \]

\[ m^2_{\tilde{\nu}}(100 \text{ GeV}) = M_2^2(100 \text{ GeV}) \left[ 1.71n_5^{-1} + 0.11 + (0.023n_5^{-1} + 0.02)t_M \right] \]

\[ - 0.5D, \quad (10c) \]

\[ m^2_{\tilde{u}_L}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[ 1.96n_5^{-1} + 0.31 + (-0.102n_5^{-1} + 0.037)t_M \right] \]

\[ - (0.5 - 0.66s^2_W)D, \quad (10d) \]

\[ m^2_{\tilde{d}_L}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[ 1.96n_5^{-1} + 0.31 + (-0.102n_5^{-1} + 0.037)t_M \right] \]

\[ + (0.5 - 0.66s^2_W)D, \quad (10e) \]

\[ m^2_{\tilde{u}_R}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[ 1.78n_5^{-1} + 0.30 + (-0.103n_5^{-1} + 0.035)t_M \right] \]

\[ - 0.66s^2_W D, \quad (10f) \]

\[ m^2_{\tilde{d}_R}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[ 1.77n_5^{-1} + 0.30 + (-0.103n_5^{-1} + 0.034)t_M \right] \]

\[ + 0.33s^2_W D, \quad (10g) \]

where \( s^2_W \equiv \sin^2 \theta_W \) and \( D \equiv -M_Z^2 \cos 2\beta \). This sfermion mass spectrum may look like that in mSUGRA in the limit when \( m_0 \ll M_{1/2} \). But that limit in mSUGRA is ruled out by the required absence of charge and colour violating vacua, as will be pointed out later. Thus the contents of the sfermion mass spectrum, specifically the squark to slepton and singlet to doublet sfermion mass ratios, distinguish mGMSB. A final point on scalar masses is that the magnitude of the \(|\mu|\) parameter is forced to become large by the requirement of EW symmetry breakdown:

\[ |\mu| \geq \frac{2}{3} n_5^{-1} M_3(M_m). \quad (11) \]

Such a large \(|\mu|\) makes the CP even charged (heavy neutral) Higgs \( H^\pm(H) \) as well as the CP odd neutral Higgs \( A \) very heavy and tightens the upper bound of 132 GeV on \( h \) in general MSSM to

\[ m_h < 120 \text{ GeV}. \quad (12) \]

The gravitino mass is given by

\[ m_{3/2} = \sqrt{\frac{1}{3}} \frac{|\langle F_3 \rangle|}{M_{PL}} = O(\text{keV}). \]
Thus the gravitino behaves here like an ultralight pseudo-Goldstino and is the LSP. If $\tilde{\chi}_1^0$ is the NLSP, it will have decays like $\tilde{\chi}_1^0 \rightarrow \gamma \tilde{G}, Z \tilde{G}, h \tilde{G}$ etc. One can estimate that

$$\tau_{\text{NLSP}} \geq 6 \times 10^{-14} \left( \frac{100 \text{ GeV}}{M_{\tilde{\chi}_1^0}} \right)^5 \left[ \frac{\Lambda_m}{(64 \lambda \text{ TeV})^2} \right]^2 \text{secs} \quad (13)$$

and $c_{\tau_{\text{NLSP}}}$ will be less than the length dimension of a detector if $M_m > 50 \text{ TeV}$. The decay photon for the $\gamma \tilde{G}$ final state provides a characteristic signature. Another interesting possibility is that of $\tilde{\tau}_1$ being the NLSP in which case one will have the prompt decay $\tilde{\tau}_1 \rightarrow \tilde{G} \tau$ and a hard, isolated $\tau$ in addition to large $E_T$ and leptons and/or jets from cascades will be a distinctive GMSB signal.

The GMSB scenario suffers from a severe finetuning problem between $|\mu|$ and $|\mu B|$. Eq. (11) makes $|\mu|$ quite large. The $\mu$ parameter originates in the GMSB scenario from a term $\lambda_\mu S H_1 H_2$ in the superpotential and a VEV $\langle s \rangle$ for the scalar component of $S$, but that leads to the soft $B\mu$ term in eq. (11) also. Then consistency with eq. (11) requires $|B| > 30 \text{ TeV}$, which is rather large and bad for the finetuning aspect in the stabilization of the weak scale.

- **Gravity mediated supersymmetry breaking**

  The messengers in this scenario are the superfields of an $N = 1$ supergravity theory, coupled to matter, with the messenger mass scale being close to the Planck scale. It has two major advantages: (1) the presence of gravity in local supersymmetry is utilized establishing a connection between global and local supersymmetry; (2) the theory automatically contains operators which can transmit supersymmetry breaking from the hidden to the observable sector. There are two disadvantages, though. First, since $N = 1$ supergravity theory is not renormalizable, one has to deal with an effective theory at sub-Planckian energies vis-a-vis poorly understood Planck scale physics. In particular, naive assumptions, made to simplify the cumbersome structure of this theory, may not hold in reality. Second, there are generically large FCNC effects of the form

$$\mathcal{L}_{\text{eff}} \sim \int d^4 \phi \ h M_{PL}^2 (\Sigma^+ \Sigma Z^+ Z), \quad (14)$$

$h$ being a typical Yukawa coupling strength.

**Lightning summary of $N = 1$ supergravity theory**

The general supergravity invariant action, with matter superfields $\Phi_i$, gauge superfields $V = V^a T^a$ and corresponding spinorial field-strength superfields $W^a$, is

$$S = \int d^6Z \left[ -\frac{1}{8} \mathcal{D} \mathcal{D} \mathcal{K} \{ \Phi^i e^V \} \Phi_j + \mathcal{W}(\Phi_i) + \frac{1}{4} f_{ab}(\Phi_i) W^a W^b \right] + h.c. \quad (15)$$

Here $\mathcal{W}$ is the superpotential, $f_{ab}(\Phi_i)$ an unknown analytic function of $\Phi$ and $\mathcal{K}$ an unknown hermitian function. The definition

$$\mathcal{G} \equiv M_{PL}^2 \left[ -3 \ln \left\{ -\frac{1}{3} M_{PL}^{-2} \mathcal{K}(\Phi^i e^V, \Phi) \right\} - \ln \left\{ M_{PL}^0 |\mathcal{W}(\Phi)|^2 \right\} \right] \quad (16)$$
and Weyl rescaling [1, 4] enable us to rewrite the non-KE terms in the integrand of eq.(15) as the potential

\[
V = -F_i G_j^F j - 3M_{PL}^4 e^{-\frac{G}{2M_{PL}^2}} + \frac{1}{2} \sum_\alpha g^2_\alpha D^{aa} D^{\alpha a},
\]

with

\[
F_i = M_{PL} e^{-\frac{G}{2M_{PL}^2}} (G^{-1})_i^j G_j + \frac{1}{4} J_{ab,k} (G^{-1})_i^j \chi^a \chi^b - (\chi^{-1})_i^k G_{kL} \chi \chi_i,
\]

\[
D^{aa} = G^i (T^{aa})_i^j \phi_j.
\]

\(G_\alpha\) being the \(\alpha\)th factor of the gauge group \(G = \prod_\alpha G_\alpha\).

The separation between the hidden sector superfields \(\Sigma\) and the observable sector ones \(Z_i\) is effected by writing

\[
\Phi_i \equiv \{Z_i, \Sigma\}, \phi_i \equiv \{z_i, \sigma\}, \Phi^i \equiv \{\bar{z}^i, \bar{\sigma}\}
\]

and assuming the additive split of the superpotential into observable and hidden parts

\[
\mathcal{W}(\Phi_i) = \mathcal{W}_0(Z_i) + \mathcal{W}_h(\Sigma).
\]

The spontaneous breakdown of supersymmetry in the hidden sector can be implemented through either a nonzero VEV \(\langle F_\Sigma \rangle\) of an auxiliary component of the \(\Sigma\) superfield or a condensate \(\langle \lambda_\Sigma \lambda_\Sigma \rangle\) of hidden sector gauginos. As a result, the gravitino becomes massive through the super-Higgs mechanism: \(m_{3/2} = M_{PL} e^{-\frac{G}{2M_{PL}^2}}\). Furthermore, soft supersymmetry breaking parameters \(A_{ijk}\) and \(B\) are generated in the observable sector with magnitudes \(\sim \langle F_\Sigma \rangle / M_{PL}\) or \(\langle \lambda_\Sigma \lambda_\Sigma \rangle / M_{PL}^2\). Scalar and gaugino masses are also generated respectively as [1, 4]

\[
m_i = O(m_{3/2}),
\]

\[
M_{ab} = \frac{1}{2} m_{3/2} \langle G^i (G^{-1})_i^j f^*_{ab,k} \rangle.
\]

The procedure suggested in Ref.[6] was to use these results as boundary conditions at the unification scale \(M_U\), where \(M_W \ll M_U < M_{PL}\), and evolve down to laboratory energies by RG equations.

**mSUGRA and beyond**

mSUGRA is a model characterized by the following specific boundary conditions on soft supersymmetry breaking parameters at the unifying scale \(M_U\):

- universal gaugino masses \(M_\alpha(M_U) = M_{1/2}, \forall \alpha\),
- universal scalar masses \(m^2_{ij}(M_U) = m^2_0 \delta_{ij}\),
- universal trilinear scalar couplings \(A_{ijk}(M_u) = A_0 \forall i, j, k\).
The soft supersymmetry breaking parameters are then treated as dynamical variables evolving from their boundary values via RG eqns. The complete set of parameters needed for mSUGRA is
\[ \{ p \} = (\text{sgn } \mu, m_0, M_{1/2}, A_0, \tan \beta). \] (21)

The magnitude $|\mu|$ of the higgsino mass gets fixed by the requirement of the EW symmetry breakdown. Among some of the immediate consequences are the predicted gaugino mass ratios at electroweak energies
\[ M_3(100 \text{ GeV}) : M_2(100 \text{ GeV}) : M_1(100 \text{ GeV}) \simeq 7 : 2 : 1 \] (22)

and the interpolating sfermion mass formulae
\begin{align*}
m_{\tilde{f}_i}^2 (100 \text{ GeV}) &= m_0^2 + 0.15M_{1/2}^2 - s_W^2 M_Z^2 \cos 2\beta, \\
m_{\tilde{f}_i}^2 (500 \text{ GeV}) &= m_0^2 + 5.6M_{1/2}^2 + (T^q_3 - Q_s s_W)M_Z^2 \cos 2\beta, \\
m_{\tilde{q}_i}^2 (500 \text{ GeV}) &= m_0^2 + 5.1M_{1/2}^2 - \frac{2}{3}s_W M_Z^2 \cos 2\beta, \\
m_{\tilde{u}_i}^2 (500 \text{ GeV}) &= m_0^2 + 5.1M_{1/2}^2 - \frac{1}{3}s_W M_Z^2 \cos 2\beta. \tag{23a-d}
\end{align*}

Let us make two final remarks on mSUGRA. First, the required absence of charge and colour violating minima disallows the limit $m_0 \ll M_{1/2}$ for mSUGRA, thereby establishing its mutual exclusivity vis-a-vis the mGMSB spectrum. Second, the $\mu$-term is somewhat less of a problem here than in GMSB since something like the Giudice-Masiero mechanism for generating it can be incorporated within this framework.

Going beyond mSUGRA, one sometimes pursues a constrained version of the MSSM, called CMSSM, where the radiative EW symmetry breakdown condition is not insisted upon. Moreover, separate universal masses are assumed at $M_U$ for fermions and Higgs bosons, since they supposedly belong to different representations of the GUT group. Now the parameter set is expanded to
\[ \{ p \}_{\text{CMSSM}} = \{ \mu, m_A, m_f, M_{1/2}, A_0, \tan \beta \}. \] (24)

Further, the spectrum plus associated phenomenology get related to but remain somewhat different from those in mSUGRA in having less predictivity. A basic criticism is the lack of justification for the still present subset of universality assumptions at $M_U$. But one is beset with severe FCNC problems if these are discarded. In particular, near mass degeneracy is needed for squarks of the first two generations and the same goes for sleptons.

There have been attempts to avoid such ad hoc universality assumptions and instead forbid FCNC through some kind of a family symmetry. A spontaneously broken $U(2)_F$, with doublets $L_a, R_a$ $(a = 1, 2)$ and singlets $L_3, R_3$, has been invoked for this purpose. The scheme works provided additional Higgs fields are introduced. Specifically, one needs ‘flavon’ fields $\phi^{ab}$ that are antisymmetric in $a, b$ and have the VEV $\langle \phi^{ab} \rangle = v\epsilon^{ab} = \begin{pmatrix} 0 & v \\ -v & 0 \end{pmatrix}$. 
Anomaly mediated supersymmetry breaking

This is a scenario in which the FCNC problem is naturally solved and yet many of the good features of usual gravity mediation are retained. It makes use of three branes, which are three dimensional stable solitonic solutions (of the field equations) existing in a bulk of higher dimensional spacetime — originally discovered in String Theory. Consider two parallel three branes, one corresponding to the observable and the other to the hidden sector. This means that all matter and gauge superfields belonging to one sector are pinned to the corresponding brane. The two branes are separated by a bulk distance $r_c \sim$ compactification radius. Only gravity propagates in the bulk. Any direct exchange between the two branes, mediated by a bulk field of mass $m$, say, will be suppressed in the amplitude by the factor $e^{-mr_c}$. (One assumes that there are no bulk fields lighter than $r_c^{-1}$). SUGRA fields, propagating in the bulk, get eliminated by the rescaling transformation $SZ \rightarrow Z$ where $S$ is a compensator left chiral superfield. However, this rescaling transformation is anomalous, giving rise to a loop induced superconformal anomaly which communicates the breaking of supersymmetry from the hidden to the observable sector. Being topological in origin, it is independent of the bulk distance $r_c$ and is also flavour blind. In consequence, there is no untowardly induction of FCNC amplitudes. One obtains one loop gaugino masses and two loop squared scalar masses as under

\begin{align}
M_\alpha &= M \frac{\beta(g_\alpha)}{g_\alpha}, \quad (25a) \\
m_i^2(Q) &= -\frac{1}{4} \left[ \beta(g_\alpha) \frac{d\gamma_i}{dg_\alpha} + \beta_\gamma \frac{\partial\gamma_i}{\partial Y} \right] m_{3/2}^2. \quad (25b)
\end{align}

Here $Y$ is a generic Yukawa coupling strength while $\gamma_i$ is the anomalous dimension of the $i$th matter superfield (N.B. $\gamma_{ij} = \gamma_i \delta_{ij}$). In addition, the trilinear A-couplings are given by

\[ A_{ijk} = -\frac{1}{2} (\gamma_i + \gamma_j + \gamma_k). \quad (26) \]

An interesting fallout of eq.(25a) is the numerical proportionality

\[ M_1(100 \text{ GeV}) : M_2(100 \text{ GeV}) : M_3(100 \text{ GeV}) = 2.8 : 1 : 7.1. \quad (27) \]
However, eq. (25b) leads to the disastrous consequence of physical sleptons becoming tachyonic since it implies $m^2_{\text{sleptons}}(M_W)<0$.

Various strategies have been attempted to evade the tachyonic slepton problem mentioned above. The simplest procedure, which defines the mAMSB model, is to add a universal dimensional constant $m^2_0$ to $m^2_i$. The manifest RG invariance of eq. (25b) is lost now and one obtains

$$m^2_i = C_i(16\pi^2)^{-2}m^2_{3/2} + m^2_0, \quad (28a)$$
$$A_{t,b,\tau} = (16\pi^2)^{-1}m^2_{3/2}h_{t,b,\tau}^{-1}\hat{\beta}_{ht,b,\tau}, \quad (28b)$$

where the $\hat{\beta}$'s and the $C_i$'s are given in Table 1. The main spectral feature in the bosino sector of this model is that the lightest neutralino $\tilde{\chi}^0_1$ and the lightest chargino $\tilde{\chi}^{\pm}_1$ are nearly mass degenerate, both being winolike, while the next higher neutralino $\tilde{\chi}^0_2$ is somewhat heavier. As a result, $\tilde{\chi}^{\pm}_1$ is longlived and can be observed if

$$180 \text{ MeV} < M_{\tilde{\chi}^{\pm}_1} - M_{\tilde{\chi}^0_1} < 1 \text{ GeV}.$$ 

The left selectron $\tilde{e}_L$ is also nearly mass degenerate with the right selectron $\tilde{e}_R$.

| $\hat{\beta}_{ht}$ | $h_t \left( -\frac{13}{15}g_1^4 - \frac{3}{2}g_2^4 + 8g_3^4 + h_t\hat{\beta}_{ht} + h_b\hat{\beta}_{hb} \right)$ |
|-----------------|----------------------------------------------------------------------------------|
| $\hat{\beta}_{hb}$ | $h_t \left( -\frac{7}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + h_t^2 + 6h_b^2 + h_\tau^2 \right)$ |
| $\hat{\beta}_{h\tau}$ | $h_\tau \left( -\frac{9}{5}g_1^2 - 3g_2^2 + 3h_b^2 + 4h_\tau^2 \right)$ |
| $C_Q$ | $-\frac{11}{50}g_1^4 - \frac{3}{2}g_2^4 + 8g_3^4 + h_t\hat{\beta}_{ht} + h_b\hat{\beta}_{hb}$ |
| $C_U$ | $-\frac{88}{25}g_1^4 + 8g_3^4 + 2h_t\hat{\beta}_{ht}$ |
| $C_D$ | $-\frac{22}{25}g_1^4 + 8g_3^4 + 2h_b\hat{\beta}_{hb}$ |
| $C_L$ | $-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + h_\tau\hat{\beta}_{h\tau}$ |
| $C_E$ | $-\frac{198}{25}g_1^4 + 2h_\tau\hat{\beta}_{h\tau}$ |
| $C_{H_2}$ | $-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3h_t\hat{\beta}_{ht}$ |
| $C_{H_1}$ | $-\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3h_b\hat{\beta}_{hb} + h_\tau\hat{\beta}_{h\tau}$ |

Table 1: Expressions for $C_i$'s and $\hat{\beta}$'s.
Gaugino mediated supersymmetry breaking

In this scenario [12], sometimes called -inoMSB, there are once again two separated parallel three branes in a higher dimensional bulk. But now only observable matter superfields are pinned to the corresponding brane, while gauge and Higgs superfields can propagate in the bulk. In this situation an interbrane gaugino or higgsino loop (cf. Fig. 4), in addition to the superconformal anomaly, can transmit supersymmetry breaking from the hidden to the observable sector. For several three branes, located in the bulk, the general decomposition of the Lagrangian is

\[ \mathcal{L}_D = \mathcal{L}_{\text{bulk}}(\Phi(x,y)) + \sum_j \delta^{(d-4)}(y - y_j) \mathcal{L}_j(\Phi(x,y), \chi_j(y)). \]  

(29)

In eq.(29) \( \Phi(x,y) \) is a typical superfield propagating in the bulk, whereas \( \chi_j(y) \) is a typical superfield localized on the \( j \)th brane. This type of a scenario does not seem to have any obvious problem. On the other hand, it has the following interesting features.

- \( M_{1/2} \sim m_{3/2} \sim |m_{H1}| \sim |m_{H2}| \sim |\mu B| \).
- Sleptons are never tachyonic.
- The \( \mu \) problem can be tackled.
- The near mass degeneracies \( M_{\tilde{\chi}_i^0} \sim M_{\tilde{\chi}_i^\pm}, \ m_{\tilde{\epsilon}_L} \sim m_{\tilde{\epsilon}_R} \) of mAMSB are lost.

A sample of sparticle masses for the given input parameters is shown in Table 2.

Braneworld supersymmetry breaking

With two separated and parallel three branes in a higher dimensional bulk, one can have more general mechanisms for the transmission of supersymmetry breaking. I just have time to mention them without going into much detail. One can have scenarios [13] using the
### Table 2: Sample points in parameter space. All masses are in GeV. In the first two points, the LSP is mostly Bino, while in the third it is a right-handed slepton.

|                | Point 1 | Point 2 | Point 3 |
|----------------|---------|---------|---------|
| **inputs:**    |         |         |         |
| $M_{1/2}$      | 200     | 400     | 400     |
| $m_{H_u}^2$    | (200)^2 | (400)^2 | (400)^2 |
| $m_{H_d}^2$    | (300)^2 | (600)^2 | (400)^2 |
| $\mu$          | 370     | 755     | 725     |
| $B$            | 315     | 635     | 510     |
| $y_t$          | 0.8     | 0.8     | 0.8     |
| **neutralinos:** |        |         |         |
| $M_{\chi_1^0}$ | 78      | 165     | 165     |
| $M_{\chi_2^0}$ | 140     | 315     | 315     |
| $M_{\chi_3^0}$ | 320     | 650     | 630     |
| $M_{\chi_4^0}$ | 360     | 670     | 650     |
| **charginos:** |         |         |         |
| $M_{\chi_1^\pm}$ | 140  | 315     | 315     |
| $M_{\chi_2^\pm}$ | 350  | 670     | 645     |
| **Higgs:**     |         |         |         |
| $\tan \beta$  | 2.5     | 2.5     | 2.5     |
| $m_{h^0}$      | 90      | 100     | 100     |
| $m_{H^0}$      | 490     | 995     | 860     |
| $m_{A}$        | 490     | 1000    | 860     |
| $m_{H^\pm}$   | 495     | 1000    | 860     |
| **sleptons:**  |         |         |         |
| $m_{\tilde{e}_R}$ | 105 | 200     | 160     |
| $m_{\tilde{e}_L}$ | 140 | 275     | 285     |
| $m_{\tilde{\nu}_L}$ | 125 | 265     | 280     |
| **stops:**     |         |         |         |
| $m_{\tilde{t}_1}$ | 350 | 685     | 690     |
| $m_{\tilde{t}_2}$ | 470 | 875     | 875     |
| **other squarks:** |  |         |         |
| $m_{\tilde{u}_L}$ | 470 | 945     | 945     |
| $m_{\tilde{u}_R}$ | 450 | 905     | 910     |
| $m_{\tilde{d}_L}$ | 475 | 950     | 945     |
| $m_{\tilde{d}_R}$ | 455 | 910     | 905     |
| **gluino:**    |         |         |         |
| $M_3$          | 520     | 1000    | 1050    |
| **other parameters:** |     |         |         |
| $M_{1/2}$      | 16      | 50      | 50      |
| $\mu$          | 19      | 78      | 78      |

Randall-Sundrum ‘warped’ metric $ds^2 = e^{-2k|\rho|} dx^\mu dx^\nu \eta_{\mu\nu} + dr^2$, with $k$ real and positive, leading to a VEV $\langle W \rangle$ of the superpotential. Alternatively, one could have compactifications analogous to string compactifications on the orbifold $S^1/Z_2 \times Z_2'$. A third possibility is to study general string or Horava-Witten compactifications of M-theory, yielding two separated three branes in a bulk. The last approach seems to provide some rationale for R-parity conservation. Generically, though, these scenarios do not yield the kind of Kähler potentials required for AMSB or -inoMSB models. The other phenomenologically interesting approach, based on string compactifications, is where SUSY breaking gets mediated by dilatino fields or superpartners of moduli fields and develops gravity mediated type of a pattern at lower energies.
Conclusion

We can summarize our conclusions in four points. (1) Gauge mediated supersymmetry breaking has a distinct $\gamma(l) + E_T$ signal, but suffers from a severe $\mu$ vs $\mu_B$ problem. (2) Gravity mediated supersymmetry breaking can generate the archetypal MSSM at electroweak energies, but has generic FCNC problems requiring additional input assumptions; with an extra singlet the $\mu$ problem can be solved by the Giudice-Masiero mechanism. (3) AMSB has the advantages of the gravity mediated scenario, but no FCNC problem; solutions to the tachyonic slepton disaster tend to be ad hoc. (4) Gaugino/higgsino mediation can lead to a phenomenologically viable model, free of many of the previous problems, but the required braneworld scenario does not seem easily derivable from String Theory.

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