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Mechanisms of supersymmetry breaking in the minimal supersymmetric standard model

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Abstract. We provide a bird's eyeview of current ideas on supersymmetry breaking mechanisms in the MSSM. The essentials of gauge, gravity, anomaly and gaugino/higgsino mediation mechanisms are covered briefly and the phenomenology of the associated models is touched upon. A few statement are also made on braneworld supersymmetry breaking.

Keywords. Supersymmetry breaking; MSSM; mechanisms.

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1. Preliminary remarks

This will be a somewhat theoretical review of models and mechanisms for generating soft explicit supersymmetry breaking terms in the MSSM. There will not 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. anomaly mediated supersymmetry breaking (AMSB), gaugino mediation as well as braneworld scenarios, connect more naturally with the latter.

Our Lagrangian can be decomposed [1] as

$$\mathscr{L}_{\text{MSSM}} = \mathscr{L}_{\text{MSSM}} + \mathscr{L}_{\text{SOFT}}, \tag{1}$$

$$-\mathscr{L}_{\text{SOFT}} = \frac{1}{2}(M_1\tilde{\lambda}_0\tilde{\lambda}_0 + M_2\vec{\tilde{\lambda}}\cdot\vec{\tilde{\lambda}} + M_3\tilde{g}^a\tilde{g}^a + \text{h.c.}) + V_{\text{SOFT}}^{\text{SCALAR}}, \tag{2}$$

$$V_{\text{SOFT}}^{\text{SCALAR}} = \sum_{\tilde{f}} \tilde{f}_i^*(\mathscr{M}_{\tilde{f}}^2)_{ij}\tilde{f}_j + (m_1^2+\mu^2)|h_1|^2 + (m_2^2+\mu^2)|h_2|^2$$
$$+(B\mu h_1\cdot h_2 + \text{h.c.}) + \text{trilinear } A \text{ terms}. \tag{3}$$

The sfermion summation in (3) 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

$$m_{\text{h}} < 132 \text{ GeV},$$

which is a 'killing' prediction of the MSSM.

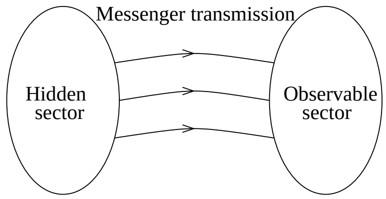

Figure 1. The transmission of supersymmetry breaking.

Though $\mathscr{L}_{\rm SOFT}$ 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 sum rule [1]:

$$S\,{\rm Tr}\,M^2_{\ell_i} + S\,{\rm Tr}\,M^2_{\nu_i} = 0 = S\,{\rm Tr}\,M^2_{u_i} + S\,{\rm Tr}\,M^2_{d_i}, \tag{4}$$

where $S\,{\rm Tr}\,M^2_f \equiv m^2_{\tilde{f}_L} + m^2_{\tilde{f}_R} - 2m^2_f$ in terms of physical masses and i is a generation index. Evidently, (4) 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 Σ which are singlets under standard model (SM) gauge transformations. Let spontaneous supersymmetry breaking (SSB) take place at a scale Λ_S in this hidden sector and be communicated to the observable world of superfields Z by a set of messenger superfields Φ (figure 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 $\sim M_s = \Lambda_S^2 M_m^{-1}$. The messengers could all be at the Planck scale (i.e. $M_m = M_{\rm Pl}$), but that need not be the case. They may or may not have nontrivial transformation properties under the SM gauge group. There 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 $M_{\rm Pl}^{-1}$.

2. Gauge mediated supersymmetry breaking [2–4]

The messenger superfields here have all the MSSM gauge interactions. MSSM superfields, with identical gauge interactions but different flavors, 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 figures 2a and 2b with $\{\phi,\chi\}$ and $\{Z,\psi\}$ being components of Φ and Z respectively. Let S be a generic hidden sector chiral superfield and $\{\Phi_i,\bar{\Phi}_i\}$ a set messenger chiral superfields [4a], interacting via couplings λ_i in the superpotential

$$W_{\rm mess} = \sum_i \lambda_i S \Phi_i \bar{\Phi}_i. \tag{5}$$

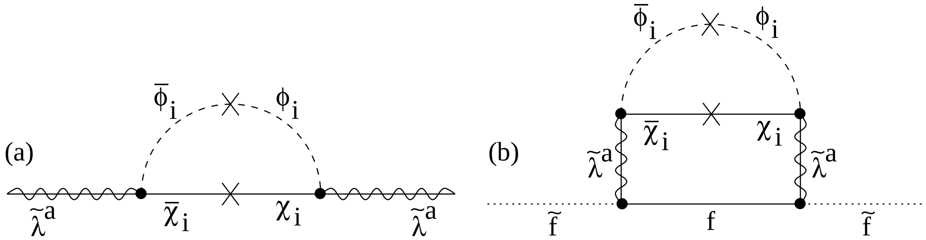

Figure 2. The origin of (**a**) gaugino and (**b**) scalar masses in GMSB.

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

$$x_i \equiv \left|\frac{\langle F_S\rangle}{\lambda_i\langle S\rangle}\right|, \qquad \Lambda \equiv \frac{|\langle F_S\rangle|}{|\langle S\rangle|}, \tag{6}$$

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 $0 < x_i < 1$.

Gaugino (scalar) masses originate in one (two) loop(s) in the manner of figure 2a(b):

$$M_\alpha = (g_2^2/16\pi^2)\Lambda\sum_\alpha 2T_\alpha(R_i)g(x_i), \tag{7}$$

$$m^2_{\tilde{f},h} = 2\Lambda^2\sum_\alpha (g_2^2/16\pi^2)^2 C_\alpha \sum_i 2T_\alpha(R_i) f(x_i). \tag{8}$$

Here Tr $T^a(\phi_i)T^b(\phi_i) = T_\alpha(R_i)\delta^{ab}$ where the trace is over the representation R_i of ϕ_i in the gauge group factor G_α and C_α is the quadratic Casimir $(\sum_a T^aT^a)_{G_\alpha}$ of the latter. Moreover,

$$g(x) = x^{-2}\left[(1+x)\ln(1+x) - (1-x)\ln(1-x)\right], \tag{9}$$

$$f(x) = x^{-2}(1+x)\left[\ln(1+x) - 2\mathrm{Li}_2\left(\frac{x}{1+x}\right) + \frac{1}{2}\mathrm{Li}_2\left(\frac{2x}{1+x}\right)\right] + (x \leftrightarrow -x), \tag{10}$$

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

$$M_\alpha \simeq (g_\alpha^2/16\pi^2)n_5\Lambda \tag{11}$$

$$m^2_{\tilde{f},h}(M_m) \simeq 2n_5^{-1}\left[C_3M_3^2(M_m) + C_2M_2^2(M_m) + \frac{3}{5}\left(\frac{Y}{2}\right)^2 M_1^2(M_m)\right], \tag{12}$$

where $C_3 = \frac{4}{3}$ (0) for an $SU(3)_C$ triplet (singlet) and $C_2 = \frac{3}{4}$ (0) for an $SU(2)_L$ doublet (singlet). To one loop, the gaugino masses (11) vary with RG evolution in the same way

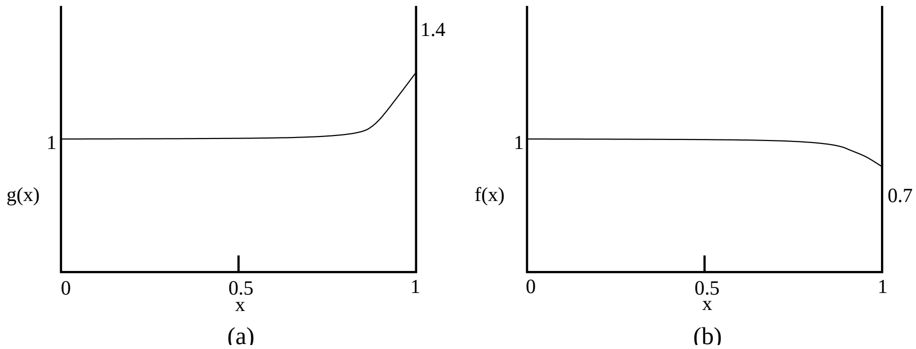

Figure 3. The behavior of (**a**) $g(x)$ and (**b**) $f(x)$.

as g_α^2, while the scalar masses (12) 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 μ, B are kept free to implement the radiative electroweak (EW) breakdown mechanism, the validity of which implies the bounds [4]

$$50 \text{ TeV} < M_m < \sqrt{n_5} \times 10^{14} \text{ GeV}. \tag{13}$$

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

$$\{p\} = \{\Lambda, M_m, \tan\beta, n_5, sgn\mu\}. \tag{14}$$

Linear RG interpolation of sfermion squarel masses from the boundary values of (12) at the scale M_m to lower energies $\sim \Lambda$ yield, with $t_M = \ln M_m/\Lambda$, the one loop expressions

$$m^2_{\tilde{e}_R}(100 \text{ GeV}) = M_1^2(100 \text{ GeV}) \times \left[1.54 n_5^{-1} + 0.05 + (0.072 n_5^{-1} + 0.01) t_M\right] + s_W^2 D, \tag{15}$$

$$m^2_{\tilde{e}_L}(100 \text{ GeV}) = M_2^2(100 \text{ GeV}) \left[1.71 n_5^{-1} + 0.11 + (0.023 n_5^{-1} + 0.02) t_M\right] + (0.5 - s_W^2) D, \tag{16}$$

$$m^2_{\tilde{\nu}}(100 \text{ GeV}) = M_2^2(100 \text{ GeV}) \times \left[1.71 n_5^{-1} + 0.11 + (0.023 n_5^{-1} + 0.02) t_M\right] - 0.5 D, \tag{17}$$

$$m^2_{\tilde{u}_L}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[1.96 n_5^{-1} + 0.31 + (-0.102 n_5^{-1} + 0.037) t_M\right] - (0.5 - 0.66 s_W^2) D, \tag{18}$$

$$m^2_{\tilde{d}_L}(500 \text{ GeV}) = M_3^2(500 \text{ GeV}) \left[1.96 n_5^{-1} + 0.31 + (-0.102 n_5^{-1} + 0.037) t_M\right] + (0.5 - 0.66 s_W^2) D, \tag{19}$$

$$m^2_{\tilde{u}_R}(500\ \text{GeV}) = M_3^2(500\ \text{GeV})\left[1.78n_5^{-1} + 0.30 + (-0.103n_5^{-1} + 0.035)t_M\right] \\ -0.66s_W^2 D, \tag{20}$$

$$m^2_{\tilde{d}_R}(500\ \text{GeV}) = M_3^2(500\ \text{GeV})\left[1.77n_5^{-1} + 0.30 + (-0.103n_5^{-1} + 0.034)t_M\right] \\ +0.33s_W^2 D, \tag{21}$$

where $s_W^2 \equiv \sin^2\theta_W$ and $D \equiv -M_Z^2\cos 2\beta$. This sfermion mass spectrum may look like that in minimal supergravity (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 color 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:

$$|\mu| \geq \frac{2}{3}n_5^{-1}M_3(M_m). \tag{22}$$

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. Furthermore, it tightens the upper bound of 132 GeV on h in general MSSM to

$$m_h < 120\ \text{GeV}. \tag{23}$$

The gravitino mass is given by

$$m_{3/2} = \sqrt{\frac{1}{3}}\frac{|\langle F_S\rangle|}{M_{\text{Pl}}} = O(\text{keV}).$$

Thus the gravitino behaves here like an ultralight pseudo-Goldstino and is the lightest supersymmetry particle (LSP). If $\tilde{\chi}_1^0$ is the NLSP, it will have decays like $\tilde{\chi}_1^0 \to \gamma\tilde{G}, Z\tilde{G}, h\tilde{G}$ etc. One can estimate that

$$\tau_{\text{NLSP}} \geq 6\times 10^{-14}\left(\frac{100\ \text{GeV}}{M_{\tilde{\chi}_1^0}}\right)^5\left[\frac{\Lambda M_m}{(64\lambda\ \text{TeV})^2}\right]^2 \text{s} \tag{24}$$

and $c\tau_{\text{NLSP}}$ will be less than the length dimension of a detector if $M_m > 50$ 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 $\tilde{\tau}_1 \to \tilde{G}\tau$ and a hard, isolated τ in addition to large $\not{E}_T$ and leptons and/or jets from cascades. This will be a distinctive GMSB signal.

The GMSB scenario suffers from a severe finetuning problem between $|\mu|$ and $|\mu B|$. Equation (22) makes $|\mu|$ quite large. The μ parameter originates in the GMSB scenario from a term $\lambda_\mu SH_1\cdot H_2$ 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. (3) also. Then consistency with eq. (22) requires $|B| > 30$ TeV, which is rather large and bad for the finetuning aspect in the stabilization of the weak scale.

3. Gravity mediated supersymmetry breaking

The messengers in this scenario [5] 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 cumbrous structure of this theory, may not hold in reality. Second, there are generically large FCNC effects of the form

$$\mathscr{L}_{\rm eff} \sim \int d^4\theta \; h M_{\rm Pl}^{-2}(\Sigma^+ \Sigma \, Z^+ Z), \tag{25}$$

h being a typical Yukawa coupling strength.

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

The general supergravity invariant action, with matter superfields Φ_i, gauge superfields $V = V^a T^a$ and corresponding spinorial field-strength superfields W^a, is [1,5].

$$S = \int d^6 z \left[-\tfrac{1}{8} \mathscr{D}\mathscr{D}\mathscr{K}\{(\Phi^\dagger e^V)_i \cdot \Phi_j\} + \mathscr{W}(\Phi_i) + \tfrac{1}{4} f_{ab}(\Phi_i) W^{aA} W^b_A \right] + \text{h.c.} \tag{26}$$

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

$$\mathscr{G} \equiv M_{\rm Pl}^2 \left[-3 \ln\{-\tfrac{1}{3} M_{\rm Pl}^{-2} \mathscr{K}(\Phi^\dagger e^V, \Phi)\} - \ln\{M_{\rm Pl}^{-6} |\mathscr{W}(\Phi)|^2\} \right] \tag{27}$$

and Weyl rescaling [1,4] enable us to rewrite the non-KE terms in the integrand of eq. (26) as the potential

$$V = -F_i \mathscr{G}^i_j \bar{F}^j - 3M_{\rm Pl}^4 e^{-\mathscr{G}/M_{\rm Pl}^2} + \frac{1}{2}\sum_\alpha g_\alpha^2 D^{\alpha a} D^{\alpha a}, \tag{28}$$

with

$$F_i = M_{\rm Pl} e^{-\mathscr{G}/(2M_{\rm Pl}^2)} (\mathscr{G}^{-1})^j_i \mathscr{G}_j + \frac{1}{4} f^*_{ab,k} (\mathscr{G}^{-1})^k_i \bar{\lambda}^a \bar{\lambda}^b - (\mathscr{G}^{-1})^k_i \mathscr{G}^{jL}_k \chi_j \chi_i, \tag{29}$$

$$D^{\alpha a} = \mathscr{G}^i (T^{\alpha a})^j_i \phi_j, \tag{30}$$

G_α being the αth factor of the gauge group $G = \prod_\alpha G_\alpha$.

The separation between the hidden sector superfields Σ and the observable sector ones Z_i is effected by writing

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

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

$$\mathscr{W}(\Phi_i) = \mathscr{W}_0(Z_i) + \mathscr{W}_h(\Sigma). \tag{31}$$

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 Σ 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: $m_{3/2} = M_{\mathrm{Pl}} \mathrm{e}^{-<\mathscr{G}>/(2M_{\mathrm{Pl}}^2)}$. Furthermore, soft supersymmetry breaking parameters A_{ijk} and B are generated in the observable sector with magnitudes $\sim \langle F_\Sigma \rangle / M_{\mathrm{Pl}}$ or $\langle \lambda_\Sigma \lambda_\Sigma \rangle / M_{\mathrm{Pl}}^2$. Scalar and gaugino masses are also generated respectively as [1,4]

$$m_i = O(m_{3/2}), \tag{32}$$

$$M_{ab} = \frac{1}{2} m_{3/2} \langle \mathscr{G}^l (\mathscr{G}^{-1})_l^k f_{ab,k}^* \rangle. \tag{33}$$

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

3.2 *mSURGA 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 $m_{ij}^2(M_U) = m_0^2 \ \delta_{ij}$,
- 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 equations. The complete set of parameters needed for mSUGRA is

$$\{p\} = (sgn\ \mu,\ m_0,\ M_{1/2},\ A_0,\ \tan\beta). \tag{34}$$

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

$$M_3(100 \text{ GeV}) : M_2(100 \text{ GeV}) : M_1(100 \text{ GeV}) \simeq 7:2:1 \tag{35}$$

and the interpolating sfermion mass formulae

$$m_{\tilde{l}_R}^2(100 \text{ GeV}) = m_0^2 + 0.15 M_{1/2}^2 - s_W^2 M_Z^2 \cos 2\beta, \tag{36}$$

$$m^2_{\tilde{l}_L}(100\ \text{GeV}) = m_0^2 + 0.53M^2_{1/2} + (T^l_{3L} - Q_l s^2_W)M^2_Z \cos 2\beta, \tag{37}$$

$$m^2_{\tilde{q}_L}(500\ \text{GeV}) = m_0^2 + 5.6M^2_{1/2} + (T^q_{3L} - Q_q s^2_W)M^2_Z \cos 2\beta, \tag{38}$$

$$m^2_{\tilde{u}_R}(500\ \text{GeV}) = m_0^2 + 5.2M^2_{1/2} + \tfrac{2}{3}s^2_W M^2_Z \cos 2\beta, \tag{39}$$

$$m^2_{\tilde{d}_R}(500\ \text{GeV}) = m_0^2 + 5.1M^2_{1/2} - \tfrac{1}{3}s^2_W M^2_Z \cos 2\beta. \tag{40}$$

Let us make two final remarks on mSUGRA. First, the required absence of charge and color violating minima disallows [7] the limit $m_0 \ll M_{1/2}$ for mSUGRA, thereby establishing its mutual exclusivity vis-a-vis the mGMSB spectrum. Second, the μ-term is somewhat less of a problem here than in GMSB since something like the Giudice–Masiero mechanism [8] 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 grand unification theory (GUT) group. Now the parameter set is expanded to

$$\{p\}_{\text{CMSSM}} = \{\mu,\ m_A,\ m_{\tilde{f}},\ M_{1/2},\ A_0,\ \tan\beta\}. \tag{41}$$

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 [9]. The scheme works provided additional Higgs fields are introduced. Specifically, one needs 'flavon' fields ϕ^{ab} that are antisymmetric in a,b and have the VEV $\langle\phi^{ab}\rangle = v\epsilon^{ab} = \begin{pmatrix} 0 & v \\ -v & 0 \end{pmatrix}$.

4. Anomaly mediated supersymmetry breaking

This is a scenario [10] 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

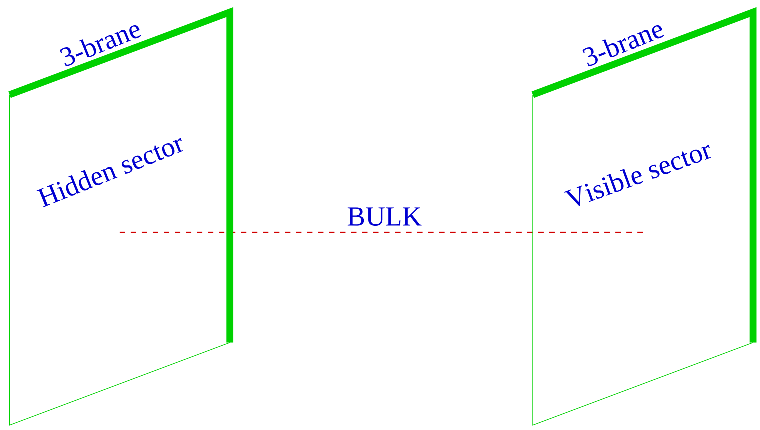

Figure 4. Hidden and observable branes in the bulk.

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 $SZ \to Z$ where S is a compensator left chiral superfield. However, this rescaling transformation is anamolous, 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 flavor 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

$$M_\alpha = M \frac{\beta(g_\alpha)}{g_\alpha}, \tag{42}$$

$$m_i^2(Q) = -\frac{1}{4}\left[\beta(g_\alpha)\frac{\mathrm{d}\gamma_i}{\mathrm{d}g_\alpha} + \beta_\gamma \frac{\partial \gamma_i}{\partial Y}\right] m_{3/2}^2. \tag{43}$$

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

$$A_{ijk} = -\tfrac{1}{2}(\gamma_i + \gamma_j + \gamma_k). \tag{44}$$

An interesting fallout of eq. (42) is the numerical proportionality

$$M_1(100\ \text{GeV}) : M_2(100\ \text{GeV}) : M_3(100\ \text{GeV}) = 2.8 : 1 : 7.1, \tag{45}$$

as contrasted with eq. (35). However, eq. (42) 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_0^2 to m_i^2. The manifest RG invariance of eq. (25b) is lost now and one obtains

Table 1. Expressions for C_i's and $\hat{\beta}$'s.

$$\hat{\beta}_{h_t} = h_t\left(-\frac{13}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + 6h_t^2 + h_b^2\right)$$
$$\hat{\beta}_{h_b} = h_t\left(-\frac{7}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + h_t^2 + 6h_b^2 + h_\tau^2\right)$$
$$\hat{\beta}_{h_\tau} = h_\tau\left(-\frac{9}{5}g_1^2 - 3g_2^2 + 3h_b^2 + 4h_\tau^2\right)$$
$$C_Q = -\frac{11}{50}g_1^4 - \frac{3}{2}g_2^4 + 8g_3^4 + h_t\hat{\beta}_{h_t} + h_b\hat{\beta}_{h_b}$$
$$C_{\bar{U}} = -\frac{88}{25}g_1^4 + 8g_3^4 + 2h_t\hat{\beta}_{h_t}$$
$$C_{\bar{D}} = -\frac{22}{25}g_1^4 + 8g_3^4 + 2h_b\hat{\beta}_{h_b}$$
$$C_L = -\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + h_\tau\hat{\beta}_{h_\tau}$$
$$C_{\bar{E}} = -\frac{198}{25}g_1^4 + 2h_\tau\hat{\beta}_{h_\tau}$$
$$C_{H_2} = -\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3h_t\hat{\beta}_{h_t}$$
$$C_{H_1} = -\frac{99}{50}g_1^4 - \frac{3}{2}g_2^4 + 3h_b\hat{\beta}_{h_b} + h_\tau\hat{\beta}_{h_\tau}$$

$$m_i^2 = C_i(16\pi^2)^{-2}m_{3/2}^2 + m_0^2, \tag{46}$$

$$A_{t,b,\tau} = (16\pi^2)^{-1}m_{3/2}h_{t,b,\tau}^{-1}\hat{\beta}_{h_{t,b,\tau}}, \tag{47}$$

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 meutralino $\tilde{\chi}_1^0$ and the lightest chargino $\tilde{\chi}_1^\pm$ are nearly mass degenerate, both being wino-like, while the next higher neutralino $\tilde{\chi}_2^0$ is somewhat heavier. As a result, $\tilde{\chi}_1^\pm$ is long-lived and can be observed [11] if

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

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

4.1 *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. figure 5), 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

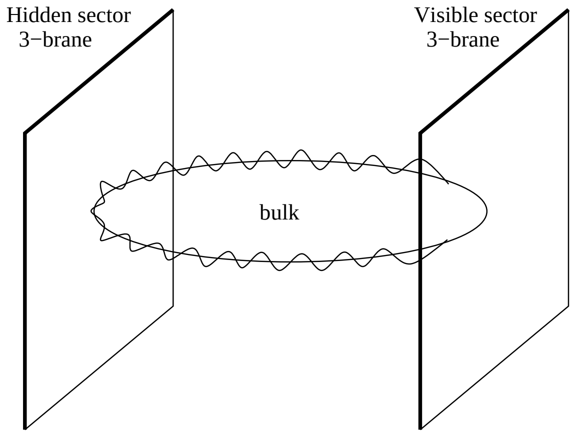

Figure 5. An interbrane -inoloop.

$$\mathscr{L}_D = \mathscr{L}_{\text{BULK}}(\Phi(x,y)) + \sum_j \delta^{(d-4)}(y-y_j)\mathscr{L}_j(\Phi(x,y),\chi_j(y)). \tag{48}$$

In eq. (48) $\Phi(x,y)$ is a typical superfield propagating in the bulk, whereas $\chi_j(y)$ is a typical superfield localized on the jth brane. This type of a scenario does not seem to have any obvious problem. On the other hand, it has the following interesting features.

- $M_{1/2} \sim m_{3/2} \sim |m_{H_1}| \sim |m_{H_2}| \sim |\mu B|$.
- Sleptons are never tachyonic.
- The μ problem can be tackled.
- The near mass degeneracies $M_{\tilde{\chi}_1^0} \sim M_{\tilde{\chi}_1^\pm}$, $m_{\tilde{e}_L} \sim m_{\tilde{e}_R}$ of mAMSB are lost.

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

4.2 *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 Randall–Sundrum 'warped' metric $\mathrm{d}s^2 = \mathrm{e}^{-2k|r|}\mathrm{d}x^\mu \mathrm{d}x^\nu \eta_{\mu\nu} + \mathrm{d}r^2$, with k real and positive, leading to a VEV $\langle \mathscr{W} \rangle$ of the superpotential. Alternatively, one could have compactifications [14] analogous to string compactifications on the orbifold $S^1/Z_2 \times Z_2'$. A third possibility [15] 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 [16], 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.

Table 2. Sample points in parameter space. All masses are in GeV. In the first two points, the LSP is mostly a Bino, while in the third it is a right-handed slepton.

		Point 1	Point 2	Point 3
Inputs	$M_{1/2}$	200	400	400
	$m^2_{H_u}$	$(200)^2$	$(400)^2$	$(400)^2$
	$m^2_{H_d}$	$(300)^2$	$(600)^2$	$(400)^2$
	μ	370	755	725
	B	315	635	510
	y_t	0.8	0.8	0.8
Neutralinos	$M_{\chi^0_1}$	78	165	165
	$M_{\chi^0_2}$	140	315	315
	$M_{\chi^0_3}$	320	650	630
	$M_{\chi^0_4}$	360	670	650
Charginos	$M_{\chi^\pm_1}$	140	315	315
	$M_{\chi^\pm_2}$	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
	μ	19	78	78

5. Conclusion

We can summarize our conclusions in four points. (1) Gauge mediated supersymmetry breaking has a distinct $\gamma(l) + \not{E}_T$ signal, but suffers from a severe μ vs. μB 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 μ 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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References

[1] M Drees, R M Godbole and P Roy, *Theory and phenomenology of Sparticles* (World Scientific, Singapore) (in press)
[2] M Dine and A E Nelson, *Phys. Rev.* **D48**, 1227 (1993)
[3] C F Kolda, *Nucl. Phys. Proc. Suppl.* **62**, 266 (1998)
[4] G F Giudice and R Rattazzi, *Phys. Rep.* **322**, 419 (1999)
[4a] Φ_i and $\bar{\Phi}_i$ together form a vector-like representation of $SU(5)$
[5] H P Nilles, *Phys. Rep.* **110**, 1 (1984)
[6] L J Hall, J Lykken and S Weinberg, *Phys. Rev.* **D27**, 2359 (1982)
[7] J A Casas, A Lleyda and C Munoz, *Nucl. Phys.* **B471**, 3 (1996)
[8] G F Giudice and A Masiero, *Phys. Lett.* **B206**, 480 (1988)
[9] R Barbieri, G Dvali and L J Hall, *Phys. Lett.* **B377**, 76 (1996)
[10] L Randall and R Sundrum, *Nucl. Phys.* **B557**, 79 (1999)
G F Giudice, M A Luty, H Murayama and R Rattazzi, *JHEP* **9812**, 27 (1998)
T Gherghetta, G F Giudice and J D Wells, *Nucl. Phys.* **B59**, 27 (1999)
[11] D Ghosh, P Roy and S Roy, *JHEP* **0008**, 031 (2000)
D Ghosh, A Kundu, P Roy and S Roy, *Phys. Rev.* **D64**, 115001 (2001)
[12] D E Kaplan, G Kribs and M Schmaltz, *Phys. Rev.* **D62**, 035010 (2000)
Z Chacko, M Luty, A Nelson and E Ponton, *JHEP* **0001**, 003 (2000)
[13] R Kallosh and A D Linde, *JHEP* **0002**, 005 (2000)
T Gherghetta and A Pomerol, *Nucl. Phys.* **B586**, 141 (2000); **B602**, 3 (2001)
E Bergshoeff, R Kallosh and A van Proeyen, *Fortschr. Phys.* **49**, 625 (2001)
[14] J Bagger, F Feruglio and F Zwirner, *Phys. Rev. Lett.* **88**, 101601 (2002)
[15] A Ansimov, M Dine, M Graesser and S Thomas, *Phys. Rev.* **D65**, 105011 (2002)
[16] S A Abel, B C Allanach, F Quevedo, L Ibanez and M Klein, *JHEP* **0012**, 026 (2000)