Phenomenology of deflected anomaly-mediation

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Abstract

We explore the phenomenology of a class of models with anomaly-mediated supersymmetry breaking. These models retain the successful flavor properties of the minimal scenario while avoiding the tachyons. The mass spectrum is predicted in terms of a few parameters. However various qualitatively different spectra are possible, often strongly different from the ones usually employed to explore capabilities of new accelerators. One stable feature is the limited spread of the spectrum, so that squarks and gluinos could be conceivably produced at TEVII. The lightest superpartner of standard particles is often a charged slepton or a neutral higgsino. It behaves as a stable particle in collider experiments but it decays at or before nucleosynthesis. We identify the experimental signatures at hadron colliders that can help distinguish this scenario from the usual ones.

1 Introduction

The origin of supersymmetry breaking is the central issue in the construction of a realistic supersymmetric extension of the Standard Model (SM). If supersymmetry is to be of any relevance to the hierarchy problem the sparticle masses should be smaller than about a TeV. Then, flavor violating processes mediated by virtual sparticles constrain their masses to preserve flavor to a high degree. One main goal of model building is to provide flavor symmetric soft terms in a simple and natural way. Gauge mediated supersymmetry breaking (GMSB) [1] represents an elegant solution to this problem: soft terms are calculable and are dominated by a flavor symmetric contribution due to gauge interactions. Supergravity, on the other hand, provides perhaps the simplest way to mediate supersymmetry breaking [2]. However, in the absence of a more fundamental theory, soft terms are not calculable in supergravity, so there is little control on their flavor structure. More technically, one could say that soft terms are dominated by "extreme ultraviolet" dynamics in supergravity and consequently are sensitive to all possible new sources of flavor violation, not just the "low-energy" Yukawa couplings. This can be considered a generic problem of soft terms mediated by supergravity. Various solutions have been suggested, including special string inspired scenarios (dilaton dominance) and horizontal symmetries.

Recently, important progress has been made in our understanding of a class of calculable quantum effects in supergravity [3, 4]. These effects can be characterized as the *pure* supergravity contribution to soft terms. This is because they are simply determined by the vacuum expectation value of the auxilliary scalar field F_ϕ in the graviton supermultiplet. The couplings of F_ϕ to the minimal supersymmetric standard model (MSSM) are a purely quantum effect dictated by the conformal anomaly. The resulting *anomaly mediated* contribution to sparticle masses is of order $\alpha F_\phi/4\pi \sim \alpha m_{3/2}/4\pi$. In a generic supergravity scenario, this calculable effect would only represent a negligible correction to the uncalculable $\sim m_{3/2}$ tree level terms. However it is consistent to consider a situation where Anomaly Mediation (AM) is the leading effect. Indeed, as pointed out by Randall and Sundrum [3], this may happen in an extra-dimensional scenario, for example, when the MSSM lives on a 3-brane, while the hidden sector lives on a brane that is far-away in a bulk where *only* gravity propagates. Recently an explicit realization of this setup has been given in ref. [5]. A more conventional situation, where the anomaly mediated contribution to just the gaugino masses and A-terms dominates, is dynamical hidden sector models without singlets [4]. Various technical aspects of AM have been further discussed in Refs. [6, 7, 8], the latter of which gives a more formal derivation along with a comparison to previous computations of quantum contributions to soft terms [9].

In pure Anomaly Mediation sfermion masses are dominated by an infrared contribution, so they are only sensitive to the sources of flavor violation that are relevant at low energy, as encoded in the fermion masses and CKM angles of the SM. Therefore AM, like the SM, satisfies natural flavor conservation. Sfermion masses are in practice family independent, since the gauge contributions dominate, like in GMSB. Unfortunately, this is not the full story: flavor is fine but the squared slepton masses are predicted to be negative.

Various attempts have been made to save the situation. In principle adding an extra supergravity contribution ruins predictivity. Nevertheless, if one assumes that some unspecified flavor universal contribution lifts the sleptons, then the low-energy phenomenology is quite peculiar [10, 11]. Other proposals involve extra fields at, or just above, the weak scale [12, 7]. In this paper we will focus on the idea of ref. [6], which we outline below.

The fact that AM provides a special Renormalization Group (RG) trajectory where all unwanted ultraviolet (UV) effects on soft terms decouple is very suggestive. Indeed, in order to solve the supersymmetric flavor problem, it would be enough to remain on this trajectory only down to a scale M_0 somewhat below the scale of flavor. In ref. [6] it was pointed out that a theory can be kicked off the AM trajectory when an intermediate theshold is governed by the vacuum expectation value (VEV) of a field X that is massless in the supersymmetric limit. This does not truly violate the UV insensitivity of AM, since the low energy theory is not just the MSSM but contains also the modulus X. While this field is coupled to the MSSM only by $1/X$ suppressed operators, its presence affects the soft masses in a relevant way. Ref. [6] used this remark to build a realistic class of models, with flavor universal and positive sfermion masses. The intermediate threshold is given by a messenger sector similar to that of GMSB models. However the sparticle spectrum of these models strongly differs from both GMSB and conventional supergravity. Indeed the prediction for gaugino mass ratios is also distinguished from "minimal" AM. The most important features of the spectrum are a reduced hierarchy between coloured sparticles and the rest, and the lightest spartner being either a slepton or a higgsino-like neutralino. The lightest supersymmetric particle (LSP) is the fermionic partner χ of the modulus X, so the lightest sparticle in the MSSM can be charged.

The purpose of the present paper is to study the implications of these novel features in collider physics and cosmology. It is organized as follows. In section 2 we recall the building blocks of the model and the corresponding high-scale boundary conditions for soft terms. In section 3 we study the low-energy spectrum and consider the constraints from electroweak symmetry breaking. In section 4 we focus on the signatures at both TEVII and LHC and draw a comparison to those of GMSB and minimal supergravity (mSUGRA). Supersymmetric corrections to rare processes are studied in section 5. In section 6 we discuss the NLSP decays and the bounds on it placed by big-bang nucleosynthesis. Section 7 contains our conclusions. In appendix A we write the one-loop RG evolution for the soft terms in terms of a minimal number of 'semi-analytic' functions, starting from the most general boundary conditions.

2 The model

Anomaly Mediated soft terms can be defined in a very simple operational way. Consider first any model in the supersymmetric limit and assign R-charge 2/3 to all its chiral matter superfields. Notice that in general this is not a true symmetry. For instance, in the superpotential only the trilinear couplings are invariant. Consider then the introduction of a spurion (classical external field) ϕ with R-charge 2/3 and scaling-dimension 1, and couple it to the original lagrangian in order to make it formally both R and scale invariant. For instance for a

generic superpotential $W(Q)$ we have

$$W(Q) = M_1 Q^2 + \lambda Q^3 + \frac{1}{M_{-1}} Q^4 + \ldots \rightarrow M_1 \phi Q^2 + \lambda Q^3 + \frac{1}{M_{-1}\phi} Q^4 + \ldots = \phi^3 W(Q/\phi). \tag{1}$$

When the choice $\phi = 1 + \theta^2 F_\phi$ is made, some special soft terms are generated: they are proportional to the dimension of the original superpotential coupling. Notice that they vanish for a purely cubic W. The same game can be played with the gauge interaction terms. Like for Yukawas, the coupling to ϕ is absent because gauge interactions are scale invariant and R symmetric at tree level. So in a theory with only gauge and Yukawa couplings no soft term arises at tree level. However, a coupling to ϕ arises at the quantum level due to anomalous breaking of scale (and R) invariance. Indeed, in order to formally restore the two symmetries one should also couple the regulator Lagrangian to ϕ. For instance in supersymmetric QED the Pauli-Villars mass should be multiplied by a factor ϕ, like in eq. (1). The quantum dependence on ϕ can be effectively accounted for by considering superfield matter wave functions and gauge couplings [13]

$$\mathcal{Z}_i(\mu) = Z\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}\right) \qquad R(\mu) = g^{-2}\left(\frac{\mu}{\sqrt{\phi\phi^\dagger}}\right) \tag{2}$$

where $Z_i(\mu)$ and $g^2(\mu)$ are the running parameters in the supersymmetric limit. Eq. 2 is derived by noticing that the quantity $\mu/\sqrt{\phi\phi^\dagger}$ is the only scale and R invariant combination of μ and ϕ [3, 4]. By eq. (2) the A-terms, scalar and gaugino masses are

$$A_{ijk}(\mu) = -\tfrac{1}{2}\left(\gamma_i(\mu) + \gamma_j(\mu) + \gamma_k(\mu)\right) F_\phi \qquad \gamma_i = \frac{d\ln Z_i}{d\ln\mu} \tag{3a}$$

$$m_i^2(\mu) = -\tfrac{1}{4}\dot{\gamma}_i(\mu)|F_\phi|^2 \qquad \dot{\gamma}_i = \frac{d\gamma_i}{d\ln\mu} \tag{3b}$$

$$m_\lambda(\mu) = \frac{\beta(g^2(\mu))}{2g^2(\mu)} F_\phi \qquad \beta = \frac{dg^2}{d\ln\mu} \tag{3c}$$

where A_{ijk} is the dimensionful scalar-Yukawa analogous to the Yukawa coupling λ_{ijk}. The pure gauge contribution to scalar masses is proportional to $-\beta(g^2)$, which is positive for asymptotically free gauge theories and negative otherwise. In the MSSM neither $SU(2)_L$ nor $U(1)_Y$ is asymptotically free. So the slepton squared masses, which are dominated by the $SU(2)_L \times U(1)_Y$ contribution, are negative and the model is ruled out.

The models constructed in ref. [6] eliminate the tachyons while preserving the successful flavor properties of AM. In these models n flavors of 'messengers' Ψ_i, $\bar{\Psi}_i$ in the $\mathbf{5} + \bar{\mathbf{5}}$ of $SU(5)$ and a singlet X are added to the MSSM fields. These fields interact via the superpotential

$$W_{\text{mess}} = \lambda_\Psi X \Psi_i \bar{\Psi}_i \tag{4}$$

so the basic structure is that of GMSB models. However it is assumed that soft terms are generated by AM already in supergravity. We are interested in a situation where X gets a large VEV so that the messengers are ultra-heavy. If $\langle X \rangle$ were fixed by supersymmetric dynamics, for example by a superpotential $W(X)$, then the relation $F_X/\langle X \rangle = F_\phi$ would hold in the presence of the spurion ϕ. The messenger supermultiplets would then be split, and upon integrating them out a gauge-mediated correction to the sparticle masses would arise. By the relation $F_X/\langle X \rangle = F_\phi$, this correction would precisely adjust the soft terms to the AM trajectory of the low-energy theory, *i.e* to the beta functions of the theory without messengers. This is just an example of the "celebrated" decoupling of heavy thresholds in AM.

However in our model, X is a flat direction in the supersymmetric limit only lifted by the effects of $F_\phi \neq 0$. The effective action along $X \neq 0$ and $\Psi, \bar{\Psi} = 0$ is determined by the running wave function $Z_X(\mu)$

$$\int d^4\theta \, Z_X\left(\sqrt{XX^\dagger/\phi\phi^\dagger}\right) XX^\dagger, \tag{5}$$

and gives the effective potential

$$V(X) = m_X^2(|X|)|X|^2 \simeq \left|\frac{F_\phi}{16\pi^2}\right|^2 n\lambda_\Psi^2(X)\left[c_\lambda\lambda_\Psi^2(X) - c_i g_i^2(X)\right]|X|^2, \tag{6}$$

where $c_\lambda, c_i > 0$, and a sum over the gauge couplings g_i of the messengers is understood. If the running mass m_X^2 is positive at large X and crosses zero at some point $X = M_0$, the potential has a stable minimum around this

point [14]. There exists a choice of parameters for which this happens: the positive Yukawa term in eq. (6) may dominate in the UV while the negative gauge contribution may balance it at a lower scale. For this mechanism to work better one may imagine the presence of a new and strongly UV free messenger gauge interaction. This is because $SU(3) \times SU(2) \times U(1)$ ends up IR free by the addition of the messengers. Around the minimum, $\mathrm{Re}(X)$ gains a mass $\sim (\alpha/4\pi)^3 F_\phi$ which could be of order a few GeV, while $\mathrm{Im}(X)$ is an axion. The crucial result, evident from eq. (5), is $F_X/X = \gamma_X(M_0)F_\phi/2$, a 1-loop quantity instead of the tree level result $F_X/X = F_\phi$ we mentioned above. Therefore, when the messengers are integrated out, their gauge-mediated contribution to sparticle masses is $\mathcal{O}(\alpha^2 F_\phi)$, which represents a negligible correction to the original $\mathcal{O}(\alpha F_\phi)$ anomaly mediated masses. Thus while the gauge beta functions are modified by eliminating the messengers, the soft terms aren't adjusted to the beta functions of the low energy theory. Below the scale M_0, the RG flow is *deflected* from the AM trajectory. That is why we call this scenario Deflected Anomaly Mediation (DAM). Practically the phenomenology of this model is that of the MSSM with boundary conditions for soft terms at scale M_0 given by AM in the MSSM plus n families of messengers. We give these boundary conditions below. Notice that the addition of messengers apparently worsens the situation in that it makes the beta functions more negative. However the gaugino masses are also changed: it is the gaugino RG contribution from M_0 to m_Z that eliminates all tachyons. An example of this behaviour for a DAM model with $n = 5$ and $M_0 = 10^{15}\,\mathrm{GeV}$ is shown in fig. 1.

The model is completed by a sector whose dynamics generate μ and $B\mu$. We remind the reader that the generation of these parameters is yet another problem of simple AM. As in GMSB, it is quite easy to obtain the right μ, but it is hard to avoid $B \sim F_\phi \gg m_{\mathrm{weak}}$. These problems are avoided in DAM by considering the addition of one singlet S coupled via the superpotential

$$\int d^2\theta \left[\lambda_H S H_{\mathrm{d}} H_{\mathrm{u}} + \frac{1}{3}\lambda_S S^3 + \frac{1}{2}\lambda_X S^2 X \right] . \tag{7}$$

Along $X \neq 0$, the field S is massive and by integrating it out the following effective operator is generated

$$\int d^4\theta \left\{ H_{\mathrm{d}} H_{\mathrm{u}} \frac{\lambda_H X^\dagger}{\lambda_X X} \tilde{Z}\left(\sqrt{XX^\dagger/\phi\phi^\dagger} \right) + \mathrm{h.c.} \right\} , \tag{8}$$

where $\tilde{Z}(\mu)$ is the running wave function mixing between X and S. Eq. 8 leads to the following expressions for μ and B at the scale M_0

$$\mu = \frac{\lambda_H}{\lambda_X}\left(\gamma_X \tilde{Z} + \dot{\tilde{Z}} \right) \frac{F_\phi^*}{2} \qquad B = \frac{\gamma_X^2 \tilde{Z} - \ddot{\tilde{Z}}}{\gamma_X \tilde{Z} + \dot{\tilde{Z}}} \frac{F_\phi}{2} \tag{9}$$

where the dots represent derivatives with respect to $\ln\mu$. Both parameters are $\sim \alpha F_\phi \sim m_{\mathrm{weak}}$. Notice that even though the effective operator eq. (8) resembles those of typical GMSB models, μ and B are the right size since F_X is a 1-loop quantity.

2.1 Predictions for the soft terms renormalized at M_0

The DAM predictions for the soft terms, renormalized at the high scale M_0, in units of $m \equiv F_\phi/(4\pi)^2$, are

$$M_i = -b_i g_i^2 m \tag{10a}$$
$$A_{RR'R''} = (c_i^R + c_i^{R'} + c_i^{R''})g_i^2 m \tag{10b}$$

for any fields $RR'R''$ except

$$A_t = A_{QUH_{\mathrm{u}}} + (\beta_t - \bar{\lambda}_H^2)m. \tag{10c}$$

The scalar masses of the fields R without significant Yukawa interactions (sleptons, d-squarks and first and second generation of u-squarks) are

$$m_R^2 = -b_i c_i^R g_i^4 m^2. \tag{10d}$$

The soft masses of Higgses and third generation Q_3 and U_3 squarks also receive significant Yukawa contributions

$$m_{H_{\mathrm{d}}}^2/m^2 = -b_i c_i^L g_i^4 + \delta \tag{10e}$$
$$m_{H_{\mathrm{u}}}^2/m^2 = -b_i c_i^L g_i^4 + \delta + \lambda_t^2(-3\beta_t + 3\bar{\lambda}_H^2) \tag{10f}$$
$$m_{U_3}^2/m^2 = -b_i c_i^U g_i^4 + \lambda_t^2(-2\beta_t + 2\bar{\lambda}_H^2) \tag{10g}$$
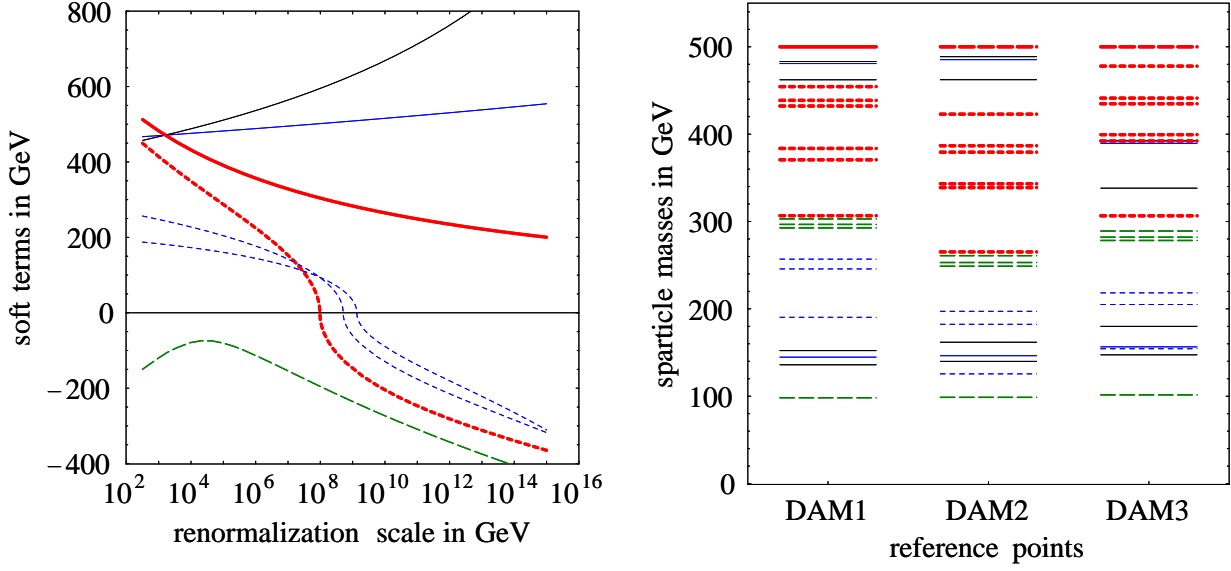$$m_{Q_3}^2/m^2 = -b_i c_i^Q g_i^4 + \lambda_t^2(-\beta_t + \bar{\lambda}_H^2) \tag{10h}$$

Figure 1: *Sample RG evolution of soft terms and qualitatively different sparticle spectra possible in DAM models. Notations are explained in the caption of fig. 4.*

where all running parameters are renormalized at M_0, $b_i = b_i^{\mathrm{MSSM}} + b_i^{\mathrm{mess}} = (33/5, 1, -3)_i + n$, the quadratic Casimir coefficients c_i^R are listed in table 4 and

$$\beta_t = (c_i^Q + c_i^L + c_i^U)g_i^2 - 6\lambda_t^2, \qquad \delta = \bar{\lambda}_H^2(4\bar{\lambda}_H^2 + 3\lambda_t^2 + \delta' - 2c_i^L g_i^2).$$

Finally, $\bar{\lambda}_H$ and δ are unknown parameters, related to the unknown parameters in the model Lagrangian as

$$\bar{\lambda}_\Psi^2 = \frac{\lambda_\Psi^2}{Z_\Psi Z_{\bar{\psi}} Z_X (1 - |\eta|^2)} \qquad \bar{\lambda}_X^2 = \frac{\lambda_X^2}{Z_S^2 Z_X (1 - |\eta|^2)^3}$$

$$\bar{\lambda}_S^2 = \frac{\lambda_S^2}{Z_S^3 (1 - |\eta|^2)^3} \qquad \bar{\lambda}_H^2 = \frac{\lambda_H^2}{Z_{H_{\mathrm{d}}} Z_{H_{\mathrm{u}}} Z_S (1 - |\eta|^2)}$$

$$\delta' = |\eta|^2 (n\bar{\lambda}_\Psi^2 + \frac{5}{2}\bar{\lambda}_X^2) + 2\bar{\lambda}_S^2 + \bar{\lambda}_X^2 - (\bar{\lambda}_S \bar{\lambda}_X^* \eta^* + \mathrm{h.c.}).$$

where $\eta = \tilde{Z}/\sqrt{Z_X Z_S}$. Notice that $|\eta| < 1$ is required for the model to be stable (positive kinetic terms). Then δ' is positive definite and δ is positive. We will see below that these extra positive contributions to the Higgs mass parameter, together with the requirement of correct electroweak symmetry breaking (EWSB), lead to an upper bound on μ/m.

As is often the case, the model-dependent couplings introduced to generate the μ and $B\mu$ terms also affect the Higgs mass parameters. In this concrete model they also affect the soft parameters of the third generation squarks. Since 4 soft masses depend on only two unknown parameters ($\bar{\lambda}_H$ and δ) there are testable predictions. On the contrary the μ and the $B\mu$ terms are determined by more than two additional unknown combinations of parameters; therefore, we consider them as free parameters and do not give their explicit expression in terms of model parameters. Even assuming real Yukawa couplings in the messenger sector, the observable sign of the $B\mu$ term is not predicted. However, if for some reason the kinetic mixing term $\tilde{Z}$ is small, CP phases can be rotated away. The model then predicts the sign of $B\mu$ and gives one relation between μ, $B\mu$ and the soft terms.

We have here neglected the effects of the other Yukawa couplings, including the possibly significantly τ and b ones. If $\tan\beta$ is large their effect should be added. They should also be taken into account when studying the predictions for 'fine details' of the spectrum (like the mass splitting between $\tilde{\tau}_1 \simeq \tilde{\tau}_R$ and $\tilde{e}_R, \tilde{\mu}_R$ and the $q/\tilde{q}$ mixing angles at the gaugino vertices induced by the CKM matrix).

The soft terms at the electroweak scale are obtained by renormalizing their values at M_0 listed in this section with the usual MSSM RG equations. The standard semi-analytic solutions cannot be applied in this case since gaugino masses do not obey unification relations, $M_i \propto \alpha_i$. In appendix A we write the RG evolution for the soft terms starting from the most general boundary conditions in terms of a minimal number of 'semi-analytic' functions. DAM models predict $M_i \propto (b_i^{\mathrm{MSSM}} + n)\alpha_i$. In this particular case the semi-analytic solutions could be further simplified.
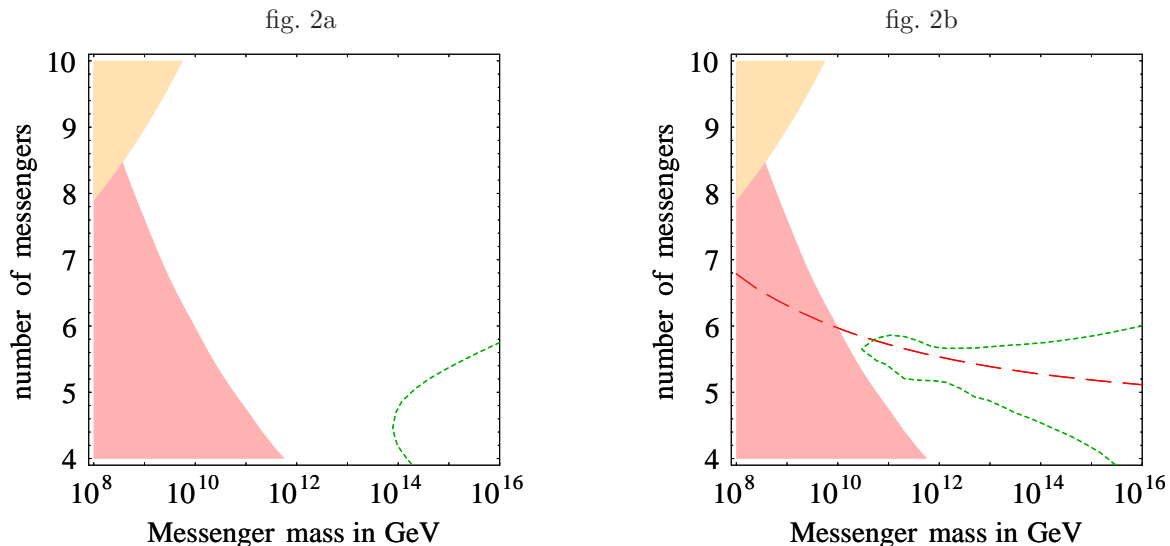
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Figure 2: *Allowed values of the main unknown model parameters, n and M_0 for $\tan\beta = 4$, $\lambda_t(M_{\text{GUT}}) = 0.5$, and small ($\bar{\lambda}_H = 0$, fig. 2a) or significant ($\bar{\lambda}_H = 1$, fig. 2a) Yukawa messengers. In the unshaded regions of the (n, M_0) plane tachyonic sleptons are avoided without too many light messengers. Below the dashed line $m^2_{H_u}$ is positive, so that EWSB is possible only with appropriate correlation between the parameters. Inside (outside) the dotted lines the lightest superpartner is a higgsino (almost always a slepton).*

3 The sparticle spectrum

The predictions for the soft terms depend on 7 parameters. The gaugino masses depend only on n (the messenger contribution to the gauge β functions); soft terms of first and second generation sfermions depend only on n and M_0 (the messenger mass); while soft terms of third generation sfermions and higgses depend also on the (imprecisely known) top Yukawa coupling at M_0, and on the possible messenger couplings $\bar{\lambda}_H$ and δ. The μ and $B\mu$ terms can be considered as free parameters, and are fixed in our analysis by the conditions of successful EWSB.

The dependence on λ_t is stronger than in gauge mediation or supergravity models. The unknown parameters $\bar{\lambda}_H$ and δ can give important corrections when n is not too large ($n \lesssim 10$): in these cases they always increase the value of $m^2_{H_u}(Q)$, and thus reduce the value of μ that gives a correct EWSB.

Even if all parameters are important, M_0 and n are the ones that control most of the sparticle spectrum (the gauginos and the sfermions). In fig. 2 we show the phenomenologically acceptable range of (M_0, n) for $\lambda_t(M_{\text{GUT}}) = 0.5$ and small messenger couplings. Shaded regions are excluded because the gauge couplings run to infinity before the unification scale (if n is too large), or because one slepton is tachyonic (if n is too low). If $n < 4$ there are tachyonic sleptons, as in pure AM where $n = 0$. If $n = 4$ sleptons can have positive squared masses, but also $m^2_{H_u}$ is positive. When $n > 4$ it is possible to have negative $m^2_{H_u}$ and positive sfermion masses unless M_0 is too low. In all the parameter space there exist unphysical deeper minima (since $m^2_\ell < 0$ at high field values, see fig. 1). There is no reason for excluding the model for this reason. Quantum and thermal tunneling rates are negligible [15]. Moreover within standard cosmology there exist plausible mechanisms [16] that naturally single out the desired physical minimum closer to the origin. A possible source of cosmological problems is the modulus X, since its mass cannot exceed a few GeV. Therefore to avoid large modulus fluctuations we must assume X to be already around its minimum when the temperature of the universe is somewhat below $\langle X \rangle$. Then, since X is only coupled to the MSSM by non-renormalizable interactions at low energy, thermal fluctuations will not affect it.

3.1 EWSB and naturalness

In most of the acceptable parameter space only the higgs field H_u has a negative squared mass term, $m^2_{H_u} < 0$, so that EWSB is induced by supersymmetry breaking in the usual way. However, $m^2_{H_u}$ is positive for certain values of the parameters: this happens for $n = 4$ (unless λ_t and $\bar{\lambda}_H$ are small); it also happens for higher values of n below the dashed lines in fig 2 if $\bar{\lambda}_H \sim 1$. With a positive $m^2_{H_u}$ it is still possible to break electroweak symmetry, but only in the narrow region of the parameter space where the μ and $B\mu$ terms give appropriate
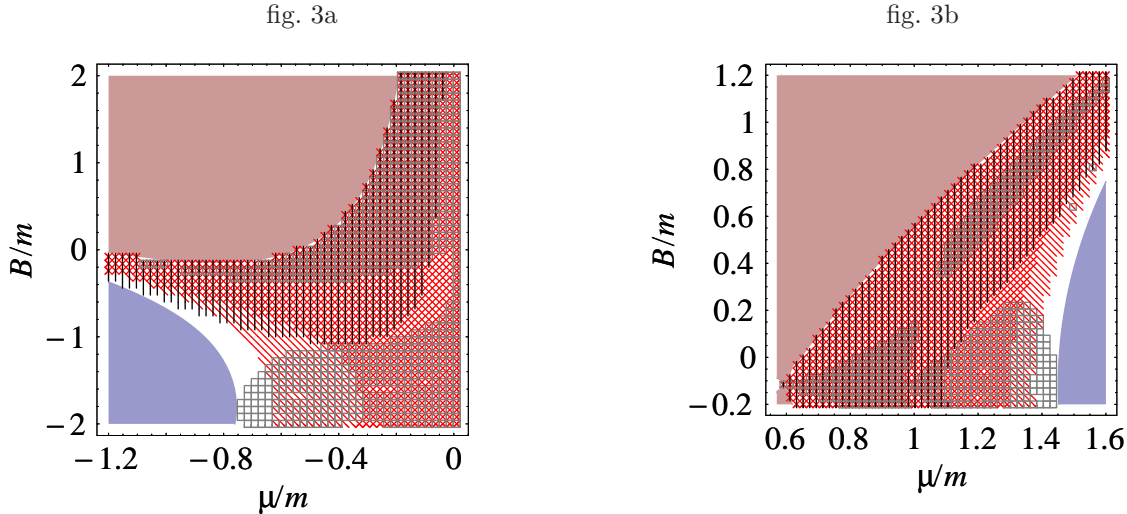
Figure 3: *The size of the allowed regions (empty regions) of the parameter space $(\mu/m, B/m)$ indicates how 'natural' is the model. Fig. 3a refers to our reference 'DAM1' model while in fig. 3b we show for comparison a mSUGRA model with $m_0 = m_{1/2}$ and $A_0 = 0$. The shaded regions are excluded because correct EWSB is not possible, while regions marked with different symbols are now experimentally excluded (see text).*

mixings in the higgs mass matrix. Moreover this situation tends to give values of $\tan\beta$ close to 1, so that the lightest higgs mass is below its experimental bound unless the sparticles are very heavy. For these reasons we do not consider this possibility attractive, and we will restrict our analysis to the more interesting case $n \geq 5$.

Strong, non-preliminary constraints on the parameter space are now given by LEP and Tevatron experiments. The bounds $m_{\chi^\pm} \gtrsim 90\,\text{GeV}$, $m_h \gtrsim 85\,\text{GeV}$ and $M_3 \gtrsim (180 \div 250)\,\text{GeV}$ are satisfied only in a small portion of the parameter space of 'conventional' supersymmetric models (like mSUGRA and GMSB), implying that the EWSB scale is unexpectedly smaller than the unobserved sparticle masses. How unnatural this situation is in any given model depends on two different characteristics of the model:

1. How light is the Z boson mass with respect to the soft terms? Since EWSB is induced by supersymmetry breaking, M_Z^2 is predicted to be a sum of various squared soft mass terms (often dominated by the gluino contribution).

2. How strong are the bounds on model parameters induced by the experimental bounds on sparticle masses? The naturalness problem becomes more stringent in the presence of an indirect bound on M_3 stronger than the direct Tevatron bound on M_3.

Concerning the second point, in SUGRA and GMSB gaugino masses obey unification relations so that the LEP bound on the chargino mass gives an indirect bound on the gluino mass, $M_3 \gtrsim 300\,\text{GeV}$, somewhat stronger than the direct Tevatron bound, $M_3 \gtrsim 220\,\text{GeV}$ (valid if $m_{\tilde{q}} \approx M_3$, as in our model). This undesired feature is not present in the scenario under study, basically for all appealing values of the parameters. However, as it happens in GMSB, the bound on the selectron mass gives an indirect bound on M_3 which is stronger than the Tevatron bound. In conclusion, for what concerns point 2, DAM is not better than 'conventional' models.

On the contrary DAM makes a somewhat more favourable prediction regarding point 1. It predicts a cancellation in the EWSB conditions for M_Z^2, because the positive radiative $\mathcal{O}(M_3^2)$ contribution to M_Z^2 is partially canceled by negative radiative $\mathcal{O}(m_{\tilde{q}}^2)$ contributions (in DAM models all sfermion squared masses are negative, before including RG corrections).

Putting it all together, DAM models suffer from some naturalness problem. This is mainly because the experimental bounds on sparticle masses are satisfied only in a small region of parameter space [17]. This is shown in fig. 3a, where we display the allowed portion of the parameter space for fixed $n = 5$, $M_0 = 10^{15}\,\text{GeV}$ and $\lambda_t(M_{\text{GUT}}) = 0.5$ and assuming that the Yukawa couplings of the messengers are negligible. With this assumption the soft terms only depend on 3 parameters: m (the overall scale of anomaly mediated soft terms), the μ-term and B. The EWSB condition allows to compute the overall SUSY scale m and $\tan\beta$ in terms of two dimensionless ratios (μ/m and B/m in figs. 3, all renormalized at M_0).

In fig. 3 we have shaded the regions where correct EWSB is not possible, and marked with different symbols the points of the parameter space where some sparticle is too light. Sampling points marked with a ($\searonearrow$,
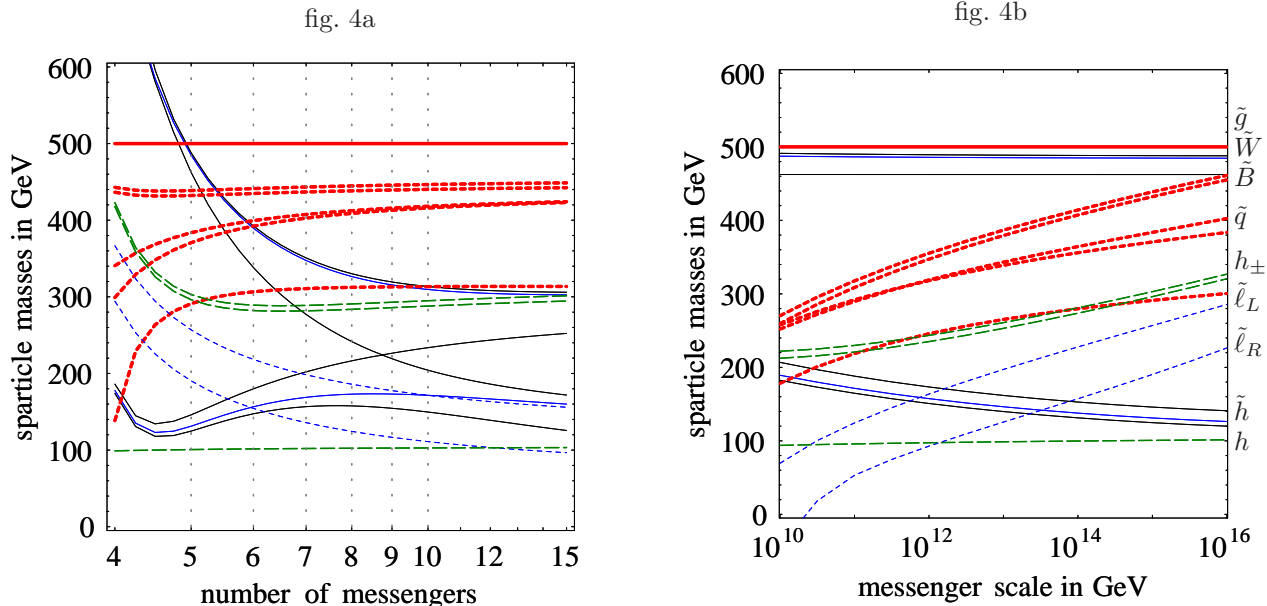
Figure 4: *Spectrum of sparticles as function of n for $M_0 = 10^{15}$ GeV (fig. 4a) and as function of M_0 for $n = 5$ (fig. 4b) at fixed $M_3 = 500$ GeV, $\lambda_t(M) = 0.5$ and negligible Yukawa messengers. Dashed (long dashed, continuous) lines refer to sfermions (higgses, fermions). Thin (thick and red) lines refer to uncoloured (coloured) sparticles. Black (blue) lines refer to the neutralinos (charginos and sleptons).*

| , —) are experimentally excluded because a (gluino, chargino, selectron, higgs) is too light. Regions where BR($B \to X_s\gamma$) differs from its SM value by more than 50% are marked with a ☐. We have restricted our plots to signs of μ and B such that the interference between charged higgs and chargino contributions to the $b \to s\gamma$ decay amplitude is destructive. With a constructive interference the indirect bounds on sparticle masses from BR($B \to X_s\gamma$) are stronger than the direct accelerator bounds and restrict the allowed parameter space to a very small region, smaller than our resolution of figs. 3.

We see that different portions of the parameter space are excluded by different combination of the bounds on gluino, charged higgsino, slepton and higgs masses. Since DAM models look somewhat disfavoured by naturalness considerations*, we also show in fig. 3b that a typical mSUGRA model (assuming $A_0 = 0$ and $m_0 = m_{1/2}$ in order to make a plot in the $(\mu/m_0, B/m_0)$ plane) has similar problems. Moreover, also gauge mediated models have a naturalness problem, mainly because they predict light right-handed sleptons. As for pure AM ($n = 0$), it predicts tachyonic sleptons, and it also has some naturalness problem: a chargino heavier than the LEP2 kinematical reach limit, $M_2 \gtrsim M_Z$, would imply that the contribution from $m_{H_u}^2$ to M_Z^2 is ~ 100 times larger than M_Z^2 itself. Adding a universal contribution to scalar masses [3, 10, 20] eliminates the tachyons but does not improve naturalness. On the other hand, DAM models also do better on the problem of naturalness.

3.2 The sparticle spectrum

We now continue our analysis studying the spectrum of sparticles in the allowed portion of the parameter space.

Before going on, we must anticipate (see the discussion in section 6) that the LSP of our models is the fermionic component of the X modulus. This fact is important as it allows a charged NSLP (sometimes a slepton). However, over the parameter space allowed in fig. 2, the NLSP decays into LSP always outside the detector. Therefore, the NLSP is practically a stable particle and the LSP plays no role in collider phenomenology.

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*In DAM models with $n \sim 5$ the fine tuning (FT [18]) of M_Z with respect to the soft terms is typically low. By choosing appropriate values of the unknown Yukawa couplings it is even possible to get FT ≈ 1. However this does not mean that DAM models are perfectly natural: since the soft terms depend on unknown Yukawa couplings the FT with respect to just the soft terms is not an adequate measure of naturalness [19].

In fig. 4 we plot the spectrum as a function of n for $M_0 = 10^{15}\,\text{GeV}$ (fig. 4a) and as a function of M_0 for $n = 5$ (fig. 4b). In both cases we have assumed $M_3 = 500\,\text{GeV}$, $\tan\beta = 4$, $\lambda_t(M_{\text{GUT}}) = 0.5$ and negligible messenger Yukawa couplings and computed the μ term from the condition of correct EWSB. Although no unique pattern emerges over all the parameter space, we try to summarize the main features of the spectrum in the following way:

0. The NLSP is usually a slepton or a neutral higgsino. The mass splitting between sleptons receives three different computable contributions; all of them (apart from a less important RG effect) tend to make an almost right-handed $\tilde{\tau}$ state the lightest slepton. Although the $\tilde{\tau}_R$ is often lighter than the higgsino (see fig. 2), it is always possible to force an higgsino NLSP by increasing the value of the unknown messenger Yukawas which decreases the value of μ that gives the correct EWSB. When $n = 4$ the NLSP can be a stop, while for large $n \gtrsim 10$ the NLSP can be a bino.

1. When $n = 5$ the NLSP is most often a neutral higgsino, sleptons are light, and all gauginos have a comparable mass *above* the squark masses.

2. When $n = 6, 7, 8$ the electroweak gauginos are lighter than the squarks, but heavier than the higgsinos.

3. When $n \gg 1$ the sfermion and gaugino masses are dominated by the pure anomaly mediated contribution to gaugino masses.

As discussed in the next section, features 1 and 2 listed above give characteristic manifestations at hadronic colliders. It is more difficult to distinguish DAM models with larger n from mSUGRA or GMSB at hadron colliders, even if for quite large values of n the mass spectrum remains significantly different from the one with unified gaugino masses. For example if $n = 20$ the ratio M_1/M_3 (connected in a simple way to the measurable ratio between the bino and the gluino masses) is still 50% higher than in the 'unified gaugino' case.

In the following section we perform more detailed studies by selecting three reference points in the DAM parameter space that capture the main characteristics of the model:

DAM1: we choose $n = 5$, $M_0 = 10^{15}\,\text{GeV}$, $\bar{\lambda}_H = 0$, $\lambda_t(M_{\text{GUT}}) = 0.5$, $M_3 = 500\,\text{GeV}$ in order to have a characteristic DAM model with $n = 5$ and higgsino NLSP.

DAM2: we choose $n = 6$, $M_0 = 10^{15}\,\text{GeV}$, $\bar{\lambda}_H = 0$, $\lambda_t(M_{\text{GUT}}) = 0.5$, $M_3 = 500\,\text{GeV}$ in order to have a characteristic DAM model with $n = 6$ and slepton NLSP.

DAM3: we choose $n = 6$, $M_0 = 10^{15}\,\text{GeV}$, $\bar{\lambda}_H = 1$, $\lambda_t(M_{\text{GUT}}) = 0.5$, $M_3 = 500\,\text{GeV}$. DAM3 is similar to DAM2, except the NLSP is a neutral higgsino.

The spectra corresponding to these three sets of parameters are shown in fig. 1 and listed in tables 1 and 2. Using these three examples we will now illustrate the phenomenology at high-energy colliders.

4 Signals at collider

The experimental manifestation of supersymmetry at hadron colliders like the Tevatron and the LHC depends strongly on how the supersymmetric particles are ordered in mass, and on the nature of the lightest superpartner of ordinary particles (stable/unstable, charged/neutral). The model under study has strong dependences on the parameters of the theory, and therefore does not make unique predictions for these important issues relevant to collider physics. Furthermore, measuring the parameters at a high-energy hadron collider is not a straightforward task. Nevertheless, we would like to point out some expectations for these models at hadron colliders despite the above difficulties.

The most important feature of the model we are presenting here is the relatively small mass gap between all the gauginos. One immediate consequence of this is a changed interpretation of gluino mass bounds from LEP2 results. The e^+e^- LEP2 collider does not produce gluinos directly, yet it does probe the Winos very effectively. Limits on the charged Wino mass from the four LEP collaborations are nearly $100\,\text{GeV}$ [21], the exact value depending on the details of the full supersymmetric spectrum. This can be interpreted as a limit on the gluino mass of about $m_{\tilde{g}} \gtrsim 300\,\text{GeV}$, provided we assume gaugino mass unification. Therefore, if the Tevatron finds a gluino with mass less than $300\,\text{GeV}$, by any of the known discovery channels, that would be one piece of evidence for the AM models. Current direct limits on the gluino mass are approximately $m_{\tilde{g}} \gtrsim 185\,\text{GeV}$ in R-parity conserving supersymmetric models with $m_{\tilde{q}} \gg m_{\tilde{g}}$, and $m_{\tilde{g}} \gtrsim 220\,\text{GeV}$ when $m_{\tilde{q}} = m_{\tilde{g}}$ [22].

Sparticle spectrum in DAM model 1

Sparticle	mass	Sparticle	mass
$\widetilde{g}$	500		
$\widetilde{\chi}_1^\pm$	145	$\widetilde{\chi}_2^\pm$	481
NLSP=$\widetilde{\chi}_1^0$	136	$\widetilde{\chi}_2^0$	152
$\widetilde{\chi}_3^0$	462	$\widetilde{\chi}_4^0$	483
$\widetilde{u}_L$	432	$\widetilde{u}_R$	384
$\widetilde{d}_L$	439	$\widetilde{d}_R$	371
$\widetilde{t}_1$	306	$\widetilde{t}_2$	454
$\widetilde{b}_1$	371	$\widetilde{b}_2$	406
$\widetilde{e}_L$	257	$\widetilde{e}_R$	190
$\widetilde{\nu}_e$	246	$\widetilde{\nu}_\tau$	246
$\widetilde{\tau}_1$	190	$\widetilde{\tau}_2$	257
h^0	98	H^0	297
A^0	293	$H^\pm$	303

Table 1: *Masses of the SUSY particles, in GeV, for the DAM model point 1.*

Sparticle spectrum in DAM model 2 | Sparticle spectrum in DAM model 3

Sparticle	mass	Sparticle	mass	Sparticle	mass	Sparticle	mass
$\widetilde{g}$	500			$\widetilde{g}$	500		
$\widetilde{\chi}_1^\pm$	176	$\widetilde{\chi}_2^\pm$	381	$\widetilde{\chi}_1^\pm$	151	$\widetilde{\chi}_2^\pm$	381
$\widetilde{\chi}_1^0$	165	$\widetilde{\chi}_2^0$	187	NLSP=$\widetilde{\chi}_1^0$	141	$\widetilde{\chi}_2^0$	162
$\widetilde{\chi}_3^0$	337	$\widetilde{\chi}_4^0$	382	$\widetilde{\chi}_3^0$	337	$\widetilde{\chi}_4^0$	382
$\widetilde{u}_L$	435	$\widetilde{u}_R$	399	$\widetilde{u}_L$	435	$\widetilde{u}_R$	399
$\widetilde{d}_L$	441	$\widetilde{d}_R$	392	$\widetilde{d}_L$	441	$\widetilde{d}_R$	392
$\widetilde{t}_1$	326	$\widetilde{t}_2$	465	$\widetilde{t}_1$	313	$\widetilde{t}_2$	470
$\widetilde{b}_1$	392	$\widetilde{b}_2$	412	$\widetilde{b}_1$	392	$\widetilde{b}_2$	410
$\widetilde{e}_L$	218	coNLSP=$\widetilde{e}_R$	154	$\widetilde{e}_L$	218	$\widetilde{e}_R$	154
$\widetilde{\nu}_e$	205	$\widetilde{\nu}_\tau$	205	$\widetilde{\nu}_e$	205	$\widetilde{\nu}_\tau$	205
coNLSP=$\widetilde{\tau}_1$	154	$\widetilde{\tau}_2$	218	$\widetilde{\tau}_1$	154	$\widetilde{\tau}_2$	218
h^0	99	H^0	283	h^0	101	H^0	290
A^0	278	$H^\pm$	289	A^0	286	$H^\pm$	296

Table 2: *Masses of the SUSY particles, in GeV, for the DAM model point 2 (left columns) and for DAM model point 3 (right columns).*

To be convinced that the DAM model is correct, much additional evidence must be gathered consistent with the model. Useful observables at hadron colliders include total rates above background in large lepton/jet multiplicity events with missing energy, invariant mass peaks of decaying heavy particles, kinematic edges to lepton or jet invariant mass spectra, and exotic signatures such as a highly ionizing track associated with a stable, heavy, charged particle track passing through the detector. All of these methods can be used to uncover evidence for supersymmetry and to help determine precisely what model is being discovered.

DAM models have several gross features that may be keys to distinguishing them from other models, such as mSUGRA and minimal GMSB. One such feature that we mentioned above is the relatively small mass difference between all the sparticles in the spectrum. Typical parameter choices in models of mSUGRA and especially GMSB have nearly an order of magnitude difference between the lightest supersymmetric partner (not counting the gravitino) and the heaviest partner. The heaviest of these sparticles are usually the strongly interacting squarks and gluinos. Consequently, unless sparticles are much heavier than the top quark, in DAM models the decays $\tilde{g} \to \tilde{t}_2 t$ and $\tilde{t}_1 \to tN$ are usually kinematically forbidden ($\tilde{t}_1$ is the lighter stop and $\tilde{t}_2$ is the heavier stop). Therefore in DAM models it is not unusual to have at most two top quarks per event, while four top quarks can be present in mSUGRA and GMSB models.

Number of leptons	DAM Model 1	DAM Model 1 with $M_i = \alpha_i M_3/\alpha_s$	DAM Model 3
0	813 (741)	714 (700)	161 (122)
1	85 (129)	105 (117)	169 (137)
2	24 (48)	12 (13)	233 (248)
3	1 (5)	1 (2)	99 (117)
4	0 (0)	1 (1)	57 (84)
5	0 (0)	0 (0)	9 (17)
6	0 (0)	0 (0)	2 (5)
7+	0 (0)	0 (0)	0 (0)

Table 3: *From DAM model point 1, the lepton multiplicity in 1000 simulated LHC events with at least $200\,\mathrm{GeV}$ of total missing energy. Leptons are counted if they have $\eta < 3$ and $p_T > 10$ GeV. The numbers in parenthesis have no p_T cut on the leptons. In the third column the spectrum is the same as DAM model point 1, except the M_1 and M_2 masses are GUT normalized. In the last column, the lepton multiplicity is given for DAM model point 3, which has significant production of leptons due to on-shell cascades of $\tilde{q} \to W \to \tilde{l}$.*

4.1 Total rates with two stable sleptons

Over much of the DAM parameter space the lightest supersymmetric partner to be produced in the detector is the $\tilde{l}_R$. For example, analyzing DAM model point 2 (see Table 2), we find that the NLSP is $m_{\tilde{l}_R} = 155\,\mathrm{GeV}$ and M_1, M_2, M_3, μ are $334, 364, 500$ and $-176\,\mathrm{GeV}$ respectively. Production of gauginos, squarks and sleptons all end up producing the lightest state $\tilde{l}_R$, which can be discovered rather easily by the detectors. The total supersymmetry production rate at the Tevatron with $\sqrt{s} = 2\,\mathrm{TeV}$ is more than 200 fb, and with several fb^{-1} expected at Tevatron runII, this choice of parameters for the model would be detected, despite superpartners not being kinematically accessible at LEP2. A careful analysis of run I data may even be able to discover or definitively rule out the parameter choices made for this example.

GMSB is another model that has a large parameter space for (quasi)-stable sleptons. If stable, charged tracks are discovered at the Tevatron, the first task will be to find the mass of the particle, and then determine the rest of the spectrum that gave rise to this sparticle. Finding the mass is relatively straightforward once there is a significant signal. Timing information along with dE/dx measurements as the particle passes through the detector are useful in this regard. Determining what model these stable tracks come from is much more difficult. One beginning step will be to estimate total supersymmetry production rate based on all $\sigma(\tilde{l}_R \tilde{l}_R + X)$ signatures. This can then be compared between the DAM model presented here and, say, minimal GMSB.

If we apply the slepton and chargino mass limits from LEP2 to GMSB, and then analyze expectations for the Tevatron, we find that squark and gluino production are not significant in supersymmetry searches at the Tevatron. This is even true when the $\tilde{l}_R$ is the NLSP and does not decay in the detector – perhaps the most likely possibility [23] in GMSB with $N_{5+\bar{5}} \geq 2$. Neglecting potentially important detector efficiency issues, every event that produces superpartners will be registered and tagged as a supersymmetry event since stable sleptons yield such an exotic signature in the detector [25, 26]. Production of sleptons, gauginos, higgsinos, and squarks all will decay ultimately to two charged sleptons plus standard model particles. Therefore, we can speak about the discovery of these models solely by analyzing the two sleptons and ignoring all other associated particles in the events, just as we did for the DAM. In this case, there is very little variability in the total rate for $\tilde{l}_R \tilde{l}_R + X$, and the rate depends mostly on the number of $5 + \bar{5}$ messengers. In Fig. 5 we plot the range allowed for total supersymmetry production [24] in GMSB with moderate to small $\tan\beta$ as a function of $m_{\tilde{e}_R}$. (Distinguishing between low and high $\tan\beta$ can be accomplished by careful analysis of the associated particles in X [26].) The upper line corresponds to $N_{5+\bar{5}} = 2$ and the lower line corresponds to $N_{5+\bar{5}} = \infty$. In contrast, recall from the paragraph above that a typical DAM set of parameters yielded a total cross-section of 200 fb for $m_{\tilde{l}_R} = 155\,\mathrm{GeV}$ because the squarks and gluinos are much lighter and contribute to the signal. Therefore, a first step in distinguishing between DAM models and GMSB models is to measure $m_{\tilde{e}_R}$ directly from stable, charged particle track analysis, and then compare the total measured rate of $\sigma(\tilde{l}_R \tilde{l}_R + X)$ to Fig. 5.

4.2 Lepton multiplicity and p_T distributions

Other important observables in supersymmetric events are the lepton multiplicity and p_T distributions. These are often sensitive to the mass hierarchies in the supersymmetric model. For example, a large source of high
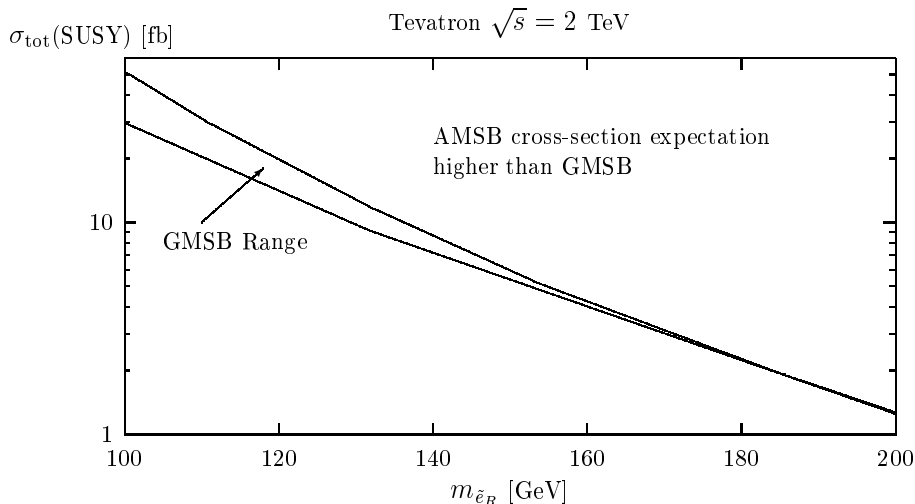
Figure 5: *Total cross-section for supersymmetry production at the Tevatron. The upper line is for GMSB with $N_{5+\bar{5}} = 2$ messengers, and the lower line is for $N_{5+\bar{5}} = \infty$ messengers. Minimal GMSB models are expected to fall within these two lines. DAM models, by contrast, are expected to be have much higher cross-sections since squark and gluinos masses are generally much lighter for the same $m_{\tilde{e}_R}$.*

p_T leptons in mSUGRA models is the cascade decays through $\chi_1^{\pm} \to l\nu\chi_1^0$. The mass difference between $m_{\chi_1^{\pm}}$ and $m_{\chi_1^0}$ is large, and the $\chi_1^{\pm}$ state is expected to participate significantly in the cascade decays of the heavier squarks and gluino down to the LSP.

In contrast, the DAM model has relatively fewer sources of high p_T leptons because of the near degeneracy of the NLSP and the next least massive chargino and neutralino. For example, in DAM model point 1 (see Table 1), we find the quasi-stable NLSP is $\chi_1^0 \sim \tilde{H}$, $M_1, M_2, M_3 = 461, 468, 500$ GeV and $380\,\text{GeV} \lesssim m_{\tilde{q}} \lesssim 440\,\text{GeV}$. Production of gluinos and squarks, while large in this model, more rarely produce sleptons because $m_{\tilde{q}} \lesssim M_1, M_2$. Instead, $\tilde{q}$ like to decay directly to a quark and a Higgsino with no intermediate leptons in a cascade decay. Leptons can arise however from $\chi_1^- \to l^- \nu\chi_1^0$, but these leptons are somewhat softer because of the near degeneracy between the mostly Higgsino χ_1^- and χ_1^0 states. In the particular example given here, the mass splitting between the lightest chargino and the lightest neutralino is about 10 GeV. The other significant source of leptons comes from third family superpartner production and decay. Since the stop and sbottom squarks are rather light in this example, many leptons do get produced from decays of the W and b particles in $t \to bW$ decays.

The lepton multiplicity and lepton p_T depend on the M_1 and M_2 masses. We illustrate this dependence by first calculating lepton observables for our example model point 1, and then doing the calculation for the same model but with M_1 and M_2 redefined to be equal to $M_i = \alpha_i M_3/\alpha_s$, consistent with gaugino mass unification, while M_3 remains the same. In this case, M_1 and M_2 are reset to 75 GeV and 145 GeV respectively, and M_3 remains at 500 GeV. In Table 4 we list the total multiplicities of leptons in 1000 simulated LHC events for the DAM example model, and the DAM example model with M_1 and M_2 redefined. The lepton multiplicity is defined to be the number of charged leptons of first and second generation with pseudo-rapidity $\eta < 3$ and transverse momentum $p_T > 10$ GeV present in each supersymmetry event. We have also required the missing energy to be greater than 200 GeV in these events to reduce standard model background, and we have not counted leptons that originate in a QCD jet (isolation requirement).

The lepton multiplicity tends to higher values for the GUT normalized gaugino spectrum rather than the untampered DAM gaugino spectrum. This is largely because more leptons pass the $p_T > 10$ GeV cut due to the large mass mass gap between the mostly bino NLSP and the next higher mass chargino and neutralino. If we put no cut on the p_T of the lepton, the number of leptons from the DAM model would be larger than the number of leptons generated in the cascade decays of the GUT normalized gaugino version of the spectrum. (The number of leptons produced with arbitrarily low p_T values is listed in parenthesis in the table.) This is indicative of the importance of looking carefully at the p_T spectrum of the leptons to see the imprint of different
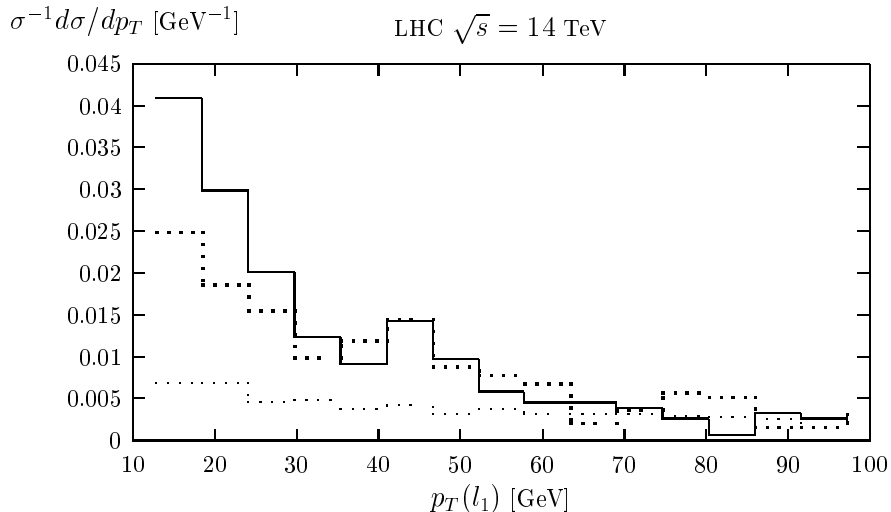
Figure 6: *The p_T distribution of the leading lepton in simulated events of supersymmetry production at the LHC. The solid line is for DAM model point 1 described in the text. The dashed line is for the same model except the electroweak gaugino masses are GUT normalized with respect to the gluino ($M_i = \alpha_i M_3/\alpha_3$). The fainter dotted line represents DAM model point 3. All lines are normalized to 1 and not the total cross-section. (The total* lepton $+ X$ *cross-section of DAM3 is a factor of 2.8 times that of DAM1.)*

mass hierarchies in the spectrum.

We demonstrate the softer lepton p_T distributions of the DAM model in Fig. 6. We have simulated 1000 supersymmetric events at the LHC and plotted the p_T distribution of the leading lepton with $\eta < 3$ and $p_T > 10\,\text{GeV}$. The effect is present as anticipated, and the magnitude of the effect is rather sizeable. In the first bin there is nearly a 50% difference between the models. We expect this observable, along with other observables [27], such as kinematic endpoint distributions, to play a key role in helping to distinguish DAM models from their competitors. In this analysis we have been assuming that the signal with large missing energy, large lepton multiplicity and large overall rate will render the standard model background not significant enough to diminish our conclusions, but of course a full investigation of the background, and simulations of real detector effects are necessary to make definitive statements about parameter determinations in supersymmetric models. Nevertheless, we are encouraged that distinctions between closely related models of supersymmetry can be made at hadron colliders.

If $n \geq 6$ it is still possible to have Higgsino NLSP: the change in the scalar masses of H_u and $\tilde{t}_i$ with $\bar{\lambda}_H \sim 1$ can alter the conditions for EWSB to allow a Higgsino NLSP. For example, if we employ the same choices of parameters that we used to generate DAM model point 2, except now we set $\bar{\lambda}_H = 1$, the resulting spectrum has a Higgsino NLSP. This is model point 3 given in the right two columns of Table 2. The phenomenology of this model with $n = 6$ and Higgsino NLSP is dramatically different than the phenomenology of point 1. In contrast to DAM1, DAM3 has a high multiplicity of leptons and high p_T distribution of leptons. Table 4 lists the lepton multiplicities for model point 3, and the faint dotted curve of Fig. 6 demonstrates the flat distribution of lepton p_T, characteristic of a high p_T spectrum of leptons. These results are readily understood by inspecting the mass hierarchies of point 3 compared to point 1. In point 3 the strongly interacting sparticles (squarks and gluinos) will almost always cascade decay to a lepton. The most effective path is through $\tilde{q} \to \tilde{W} \to \tilde{l}$, where at least one lepton results. The mass hierarchies of point 1 do not allow these high lepton multiplicities. Therefore, close inspection of the lepton observables may provide a handle on the parameter $\bar{\lambda}_H$ in addition to measurements of the various sparticle masses.

5 Signals in rare processes

In DAM models the soft terms could contain no extra flavour or CP violating terms beyond the ones induced by the CKM matrix. There are however two possible exceptions.

1. The μ and B terms could be complex: in this case they would typically generate too large electron and neutron electric dipoles, unless their phases are so small (less than about 0.01) [28] that do not significantly affect collider observables.

2. Extra Yukawa couplings not present in the SM can affect the soft terms in a way that crucially depends on how the soft terms are mediated. In supergravity heavy particles affect the soft terms, while in pure AM models soft terms are not affected by fields above the supersymmetry breaking scale. Like in GMSB, in DAM models the soft terms are not affected by interactions of fields heavier than the messenger mass M_0. The effective theory at M_0 however might not be the MSSM. For example, some of the 'right-handed neutrinos' N often introduced in order to generate the observed neutrino masses could be lighter than M_0. If they have order one Yukawa couplings $\lambda_N N L H_{\rm u}$ they imprint lepton flavour violation in the soft mass terms of left-handed sleptons $\tilde{L}$ inducing significant rates for processes like $\mu \to e\gamma$. Unlike in GMSB models, in DAM models these effects are not suppressed by a (RG-enhanced) loop factor. However for a right-handed neutrino mass $M_N \sim 10^{9 \div 11}$ GeV, optimal for leptogenesis, the Yukawa couplings λ_N must be small, $\lambda_N \lesssim 0.005$, in order to get a left-handed neutrino mass smaller than 1 eV.

If none of these exceptions is realized, in DAM models supersymmetric loop effects only give new contributions to processes already present in the SM ($b \to s\gamma$, $g-2$ of μ, ϵ_K, Δm_B, $K \to \pi\bar{\nu}\nu$ decays) but cannot give rise to new effects (like $\mu \to e\gamma$ decay, electric dipoles, contributions to K, B_d, B_s, D physics with non-CKM and/or non-SM chiral structure). Taking into account the accelerator bounds on sparticle masses, few rare processes can receive interesting contributions:

- Supersymmetric corrections can significantly enhance $\mathrm{BR}(B \to X_s\gamma)$ [29] over its SM value. For example in all the reference points studied in the previous section the $b \to s\gamma$ effective operator (with all fields and couplings renormalized at the relevant scale $Q \sim m_B$) is

$$\mathcal{H}_{\rm eff} = -[-0.29(\mathrm{SM}) - 0.08(\text{charged higgs}) \pm 0.07(\text{chargino})] V_{tb} V_{ts}^* \frac{e g_2^2}{(4\pi)^2} \frac{m_b}{2M_W^2} [(\bar{s}_L \gamma_{\mu\nu} F^{\mu\nu} b_R) + \text{h.c.}]$$

  Unless the chargino contribution compensates the charged higgs contribution (its sign depends on the relative sign between m, μ and B), $\mathrm{BR}(B \to X_s\gamma)$ is two times larger than in the SM, conflicting with experimental bounds. Even assuming a destructive interference (otherwise the sparticles must have unnaturally heavy masses) a detectable supersymmetric correction to the $B \to X_s\gamma$ branching ratio remains likely. In these models the gluino/bottom contribution is computable, and turns out to be negligible.

- Since EW gauginos are heavier than in mSUGRA or GM models, a supersymmetric correction to the anomalous magnetic moment of the μ [30], at a level detectable in forthcoming experiments [31] is rather unlikely (but not impossible).

- The supersymmetric corrections to K and B mixing [29] can be larger than in mSUGRA and GMSB models, because coloured sparticles can be lighter. With a 'reasonable' sparticle spectrum, Δm_B can be enhanced by $(20 \div 25)\%$ with respect to its SM value. Such corrections are comparable to the present theoretical uncertainties on the relevant QCD matrix elements. Larger corrections are present in small corners of the parameter space with light stops.

6   NLSP decays and nucleosynthesis

The lightest supersymmetric particle is the fermionic partner χ of the modulus X. Indeed by studying the effective action in eq. (5) one finds $m_\chi = \mathcal{O}(\alpha/4\pi)^2 F_\phi$. Therefore, unless some coupling in the messenger sector is strong, we expect m_χ to be smaller than a few GeV, so that χ is the LSP. The is a welcome fact: the LSP of our model is automatically neutral and unwanted charged relics are avoided. On the other hand, the lightest sparticle in the SM sector, the NLSP, can be charged (a right-handed slepton) as it decays into χ. Now, the effective couplings governing this decay are suppressed by inverse powers of the messenger mass and by loop factors. Indeed χ plays a role similar to that of the Goldstino in gauge mediated models. In the range of allowed M_0, the NLSP lifetime is so long that it behaves as a stable particle in collider experiments. However, lifetimes in excess of 1 sec, can dangerously affect the big-bang predictions of light element abundances. In the rest of this section we will discuss the constraints placed on M_0 by nucleosynthesis.

Let us first derive the couplings of χ to the SM particles. For a chiral matter multiplet Q the effective Lagrangian, computing loop corrections with superfield techniques, is [6]

$$\mathcal{L}_{\text{eff}} = \int d^4\theta \; Z_Q \left(\mu^2/\phi\phi^\dagger, XX^\dagger/\phi\phi^\dagger \right) QQ^\dagger \tag{11}$$

leading to a coupling

$$\mathcal{L}_{\chi q \tilde{q}} = \rho_q \frac{F_\phi}{M_0} \chi q \tilde{q}^\dagger + \text{h.c.} \quad \text{where,} \tag{12a}$$

$$\rho_q = (\partial_{\ln \mu^2} + \partial_{\ln XX^\dagger}) \partial_{\ln XX^\dagger} \ln Z_Q. \tag{12b}$$

The above expression is easily obtained by expanding $\ln Z_Q$ in powers of $\ln \phi$ and $\ln X/M_0$ and by noting that the leading contribution to $\mathcal{L}_{\chi q \tilde{q}}$ comes from second order cross terms $\propto \ln \phi \ln X/M_0$. In the case of right-handed sleptons we have

$$\rho_{e_R} = \frac{1}{8\pi^2} \frac{2n(n+33/5)}{11} \left(\alpha_1^2(M_0) - \alpha_1^2(M_Z) \right), \tag{13}$$

where we have taken $\mu = m_Z$ in $\mathcal{L}_{\text{eff}}$. Notice that the coupling from eqs. (12b), (13) is qualitatively similar to the Goldstino coupling for a gauge mediated model with $F_X/X = F_\phi$. However in gauge mediation, unlike here, $\rho_q = m_{\tilde{q}}^2 (M_0^2/F_X)^2$ by current algebra.

In the case of a higgsino NLSP the relevant term is the one generating μ

$$\mathcal{L}_{\text{eff}} = \int d^4\theta \; H_{\text{u}} H_{\text{d}} \frac{X^\dagger}{X} \tilde{Z} \left(\sqrt{XX^\dagger/\phi\phi^\dagger} \right). \tag{14}$$

As discussed in section 2, the effective μ term is equal to $(X^\dagger \tilde{Z})|_{\bar{\theta}^2}/M_0$. By writing $X = M_0 + \delta X$, it is easy to see that, at the leading order in an expansion in $1/M_0$ and α, eq. (14) leads to a superpotential coupling

$$\mathcal{L}_{\text{eff}} = -\int d^2\theta \frac{\mu}{M_0} H_{\text{u}} H_{\text{d}} \, \delta X. \tag{15}$$

Notice that the coupling of χ to the Higgs sector is stronger than that to sfermions. It is proportional to the supersymmetric mass μ (1-loop) rather than to the mass splitting $B\mu$ (2-loop). This is consistent, since χ is not the Goldstino. The most important consequence of eq. (15) is that it can mediate the decay $N_1 \to h\chi$ whenever allowed by phase space. For a higgsino-like NLSP, we have $m_{N_1} \simeq \mu$ with $N_1 \simeq \tilde{H}_\pm^0 = (\tilde{H}_{\text{d}}^0 \pm \tilde{H}_{\text{u}}^0)$, depending on the sign of μ. The width of a higgsino-like NLSP is then

$$\Gamma_{N_1 \to h\chi} = \frac{(\cos\alpha \mp \sin\alpha)^2}{64\pi} \frac{\mu^3}{M_0^2} \left(1 - \frac{\mu^2}{m_h^2} \right)^2 \tag{16}$$

corresponding to a lifetime shorter than a second over most of parameter space already for $M_0 \lesssim 10^{15}$ GeV. We conclude that nucleosynthesis does not place significant bounds on a higgsino LSP whenever $\mu > m_h$, which is almost required by experimental bounds.

Let us consider now the bounds on a stau NLSP. The coupling to χ is smaller than for the higgsino NLSP (2-loop versus 1-loop). The correspondingly longer $\tilde{\tau}$ lifetime is well approximated, as a function of $m_{\tilde{\tau}}$ and M_0, by

$$\tau_{\tilde{\tau}} = \left(\frac{M_0}{10^{13}\text{GeV}} \right)^2 \left(\frac{200\text{GeV}}{m_{\tilde{\tau}}} \right)^3 \quad \text{sec.} \tag{17}$$

This quantity is larger than 1 sec over a significant fraction of parameter space, where the $\tilde{\tau}$ decay can dangerously affect nucleosynthesis. The most stringent bounds come from decays processes involving hadronic showers. These showers break up the ambient ^{4}He into D and ^{3}He and can lead to an overabundance of the two latter elements. The showers can also overproduce ^{6}Li and ^{7}Li from "hadrosynthesis" of ^{3}He, T or ^{4}He. The decay $\tilde{\tau} \to \chi\tau$ leads to hadronic showers as the τ further decays hadronically with a large branching ratio. Using the results in ref. [32, 33], it was concluded in ref. [34] that lifetimes larger than 10^4 sec lead to unacceptable overproduction of ^{7}Li. Ref. [34] shows a careful analysis, including a computation of the relic NLSP density at nucleosynthesis, for gauge mediated models with a stau NLSP. A similarly detailed analysis is beyond the aim of the present paper, but we expect that the results of [34] can be carried over to our case. This is because the bounds do not depend very strongly on the $\tilde{\tau}$ relic density, which in our model is not going to differ drastically

from that in gauge mediation. Therefore we conclude that overproduction of ^{7}Li gives the bound $\tau_{\tilde{\tau}} < 10^4$ sec. By eq. (17) this bound roughly translates into $M_0 < 10^{14}$ GeV.

A stronger bound, forbidding decays between 10^2 and 10^4 sec can come from the deuterium abundance X_D normalized to hydrogen. However there is, at the moment a controversy in the measurement of X_D from astrophysical observation. Two values are quoted in the literature, a high one $X_D = (1.9 \pm 0.5) \times 10^{-4}$ from ref. [35] and a low one $X_D = (3.39 \pm 0.25) \times 10^{-5}$ from ref. [36]. In ref. [34] it was concluded that no further bounds are obtained when the high value of X_D is assumed. On the other hand, the low X_D value can give a stronger constraint $\tau_{\tilde{\tau}} < 10^2$ sec.

We conclude that nucleosynthesis places a significant bound on the messenger mass when the NLSP is a stau. This bound on M_0 can range between 10^{13} and 10^{15} GeV depending upon the model parameters $m_{\tilde{\tau}}$ and n and on the astrophysical input data. We stress that while the bound is not negligible, there remains a large allowed region $10^{10} < M_0 \lesssim 10^{14}$ GeV, where nucleosynthesis is fine.

7   Conclusions

We have studied the phenomenology of models where the presence of a light modulus X induces a calculable correction to anomaly mediated soft masses. This correction lifts the tachyonic sleptons while preserving the flavor universality of anomaly mediation. The resulting MSSM phenomenology is interesting and fairly distinguished from both minimal supergravity and gauge mediation. The gaugino masses are not unified, and the gluino is not much heavier than the other gauginos (see fig. 4). All sfermion masses start out negative at a scale between 10^{10} and 10^{16} GeV but are driven positive at a lower scale by the RG contribution of gaugino masses. Because of all these features the spectrum is a lot more compact than in minimal supergravity or gauge mediation so that coloured sparticles can be produced and studied at TEVII. Gauginos and squarks have comparable masses and are somewhat heavier than higgsinos and sleptons. The lightest superpartner is either a neutral higgsino or a right-handed stau, but it is only an NLSP. The LSP is the fermionic partner χ of the modulus X. The NLSP decay into χ takes place outside the detector. The rate of this decay does not conflict with the successful predictions of big bang nucleosynthesis over a significant portion of parameter space.

The signals of deflected anomaly mediation at hadron colliders are easily distinguished from the conventional ones. In the case of a charged slepton NLSP (e.g., DAM2) the signature is similar to GMSB with two or more messengers: two highly ionizing tracks in the detector. However, the total production cross section as a function of the slepton mass is much bigger in DAM than in GMSB. For instance at Tevatron it could be a factor 20 bigger. This is because for a given slepton mass, gluinos and squarks are about a factor of 2 lighter in DAM than in GMSB. So in case stable charged tracks are discovered, one can easily tell DAM from GMSB.

When the NLSP is higgsino (e.g., DAM1) the competing scenarios have usually bino LSP. Here the relevant observables are lepton multiplicity and p_T distributions in supersymmetric events. The signature of DAM depends crucially on which is the bigger between the squark and wino mass (each case can arise by proper parameter choices in DAM). Since $m_{\tilde{W}} > m_{\tilde{q}}$ for DAM1, squark production leads to a cascade with fewer high p_T leptons than in standard bino LSP scenarios. The softness of the leptons is due to the small mass splitting among the charged and neutral higgsinos produced in the cascade, while in mSUGRA and GMSB the LSP is well split from the next higher mass neutralino and chargino. Also, now the squarks often decay directly to the lightest higgsino, without producing any lepton.

On the other hand for $m_{\tilde{W}} < m_{\tilde{q}}$ (e.g., DAM3), more high p_T leptons are produced than usual. This is because squarks can decay via $\tilde{q}_L \rightarrow \tilde{W} \rightarrow \tilde{H}^0_{1,2}, \tilde{H}^+$ and $\tilde{q}_L \rightarrow \tilde{W} \rightarrow \tilde{l} \rightarrow \tilde{H}^0$. Energetic leptons are then produced in $\tilde{W}$ decays and/or the $\tilde{l}$ decays, while additional softer leptons are produced when $\tilde{H}_2$ and $\tilde{H}^+$ further decay to $\tilde{H}_1$. The lepton p_T distribution for the above cases is shown in Fig. 6 where it is compared to a standard bino LSP scenario, and lepton multiplicities are given in Table 3. Of course similar signatures are obtained in any scenario where the higgsinos are somewhat lighter than winos and bino. However, the unique mass hierarchy of the charginos, neutralinos and sleptons in DAM, as illustrated by the spectrum of DAM3, rarely occurs in GMSB or mSUGRA. To further tell DAM from these other possibilities one can resort to other observables. One additional consequence of the compact DAM spectrum is that more than 2 tops in the gluino cascade are often forbidden by phase space, whereas higher multiplicity of top quarks may exist in final states of GMSB and mSUGRA.

We conclude that DAM provides an interesting alternative to conventional soft term scenarios from both the theoretical and the phenomenological point of view. This example also provides hope that we may not have to wait for the LHC to discover superpartners: TEVII may have enough luminosity and energy.

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i	b_i	c_i^Q	c_i^U	c_i^D	c_i^L	c_i^E	c_i^{u}	c_i^{d}	c_i^{e}
1	$\frac{33}{5}$	$\frac{1}{30}$	$\frac{8}{15}$	$\frac{2}{15}$	$\frac{3}{10}$	$\frac{6}{5}$	$\frac{13}{15}$	$\frac{7}{15}$	$\frac{9}{5}$
2	1	$\frac{3}{2}$	0	0	$\frac{3}{2}$	0	3	3	3
3	-3	$\frac{8}{3}$	$\frac{8}{3}$	$\frac{8}{3}$	0	0	$\frac{16}{3}$	$\frac{16}{3}$	0

Table 4: Values of the RG coefficients in the MSSM.

A   RG evolution of soft terms with non unified gaugino masses

In this appendix we present semi-analytic solutions for the one-loop RG evolution of the soft terms in presence of the large Yukawa coupling of the top. We give the soft terms at an arbitrary energy scale Q, starting from an arbitrary scale M_0 with arbitrary gaugino masses M_{i0}, sfermion masses m_{R0}^2, A terms A_{g0}^{f}, μ-term μ_0, and B-term B_0. We do not assume unification of the gauge couplings. Here $i = \{1, 2, 3\}$ runs over the three factors of the SM gauge group, $\mathrm{f} = \mathrm{u}, \mathrm{d}, \mathrm{e}$, $g = 1, 2, 3$ is a generation index and R runs over all the scalar sparticles $(Q_g, U_g, D_g, E_g, L_g, H_{\mathrm{u}}, H_{\mathrm{d}})$. These formulæ, obtained with superfield techniques [13, 37], are significantly simpler than equivalent ones already existing in the literature [38] because they never involve double integrals over the renormalization scale. The running soft terms renormalized at an energy scale Q are

$$M_i(Q) \;=\; M_{i0}/f_i \tag{18a}$$

$$\mu(Q) \;=\; \mu_0 \cdot y^{b_1^{\mathrm{u}}} E_h \tag{18b}$$

$$B(Q) \;=\; B_0 + 2x_{i1}^L M_{i0} - b_1^{\mathrm{u}} I'/b_t \tag{18c}$$

$$A_g^{\mathrm{f}}(Q) \;=\; A_{g0}^{\mathrm{f}} + x_{1i}^{\mathrm{f}}(E) M_{i0} - b_g^{\mathrm{f}} I'(E)/b_t \tag{18d}$$

$$m_R^2(Q) \;=\; m_{R0}^2 + x_{i2}^R M_{i0}^2 - b_R^t I - \tfrac{3}{5}\tfrac{Y_R}{b_1} I_Y \tag{18e}$$

where

$$t(Q) \equiv \frac{2}{(4\pi)^2} \ln \frac{M}{Q} \qquad f_i(t(Q)) \equiv \frac{\alpha_i(M_0)}{\alpha_i(Q)}, \qquad E^\alpha(t) \equiv \prod_i f_i^{c_i^\alpha/b_i}(t), \qquad x_{i^n}^\alpha \equiv \frac{c_i^\alpha}{b_i}(1 - f_i^{-n})$$

and M is any scale. All b-factors are simple numerical coefficients: the b_i are the coefficients of the one-loop β functions, $\{b_1, b_2, b_3\} = \{33/5, 1, -3\}$. The Y_R are the hypercharges of the various fields R, normalized as $Y_E = +1$. The b_R^t coefficients vanish for all fields R except the ones involved in the top Yukawa coupling: $b_{H_{\mathrm{u}}}^t = 1/2$, $b_{\tilde{Q}_3}^t = 1/6$ and $b_{\tilde{U}_3}^t = 1/3$. The factor $I_Y = (1 - 1/f_1)\,\mathrm{Tr}[Y_R m_{R0}^2]$ takes into account a small RG effect induced by the $\mathrm{U}(1)_Y$ gauge coupling. The λ_t effects are contained in

$$I' \;=\; \rho[A_{t0} + M_{i0} X_i] \tag{19a}$$

$$I \;=\; \rho(m_{Q_30}^2 + m_{U_30}^2 + m_{H_{\mathrm{u}}0}^2) + (1-\rho)A_{t0}^2 +$$
$$\qquad -\big[(1-\rho)A_{t0} - M_{i0}\rho X_i\big]^2 + \rho M_{i0}^2 X_{i^2} + \rho M_{i0} M_{j0} X_{ij} \tag{19b}$$

where $A_{t0} = A_{30}^{\mathrm{u}}$ is the top A-term at M_0 and

$$X(M_0, Q) \equiv \int_{t(M_0)}^{t(Q)} E^{\mathrm{u}}(t)dt, \qquad X_{i^n}(M_0, Q) = \frac{\int E^{\mathrm{u}} x_{i^n}^{\mathrm{u}}\,dt}{\int E\,dt}, \qquad X_{ij}(M_0, Q) = \frac{\int E^{\mathrm{u}} x_{i1}^{\mathrm{u}} x_{i1}^{\mathrm{u}}\,dt}{\int E\,dt} \tag{20}$$

All the integrals are done in the same range as the first one. The 'semi-analytic' functions X_{i^n} are needed only for $n = 1$ and 2. In practice one has to compute numerically few functions of two variables, Q and M_0. A more efficient computer implementation is obtained rewriting the $X_{...}(M_0, Q)$ functions in terms of $1 + 3 + 9$ functions with only one argument

$$F(M_0) - F(Q) = \int_{t(M_0)}^{t(Q)} E^{\mathrm{u}}(t)dt, \qquad F_{i^n}(M_0) - F_{i^n}(Q) = \int \frac{E^{\mathrm{u}}}{f_i^n}dt, \qquad F_{ij}(M_0) - F_{ij}(Q) = \int \frac{E^{\mathrm{u}}}{f_i f_j}dt.$$

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