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Anomaly Mediation, Fayet-Iliopoulos D-terms and precision sparticle spectra

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Abstract

We consider the sparticle spectra that arise when anomaly mediation is the source of supersymmetry-breaking and the tachyonic slepton problem is solved by a Fayet-Iliopoulos D-term. We also show how this can lead to a minimal viable extension of anomaly mediation, in which the gauge symmetry associated with this D-term is broken at very high energies, leaving as its footprint in the low energy theory only the required D-terms and seesaw neutrino masses.

1 Introduction

Increasing precision of sparticle spectrum calculations is an important part of theoretical preparation for the LHC and the ILC. Much of this work has concentrated on the MSUGRA scenario, where it is assumed that the unification of gauge couplings at high energies is accompanied by a corresponding unification in both the soft supersymmetry-breaking scalar masses and the gaugino masses; and also that the cubic scalar interactions are of the same form as the Yukawa couplings and related to them by a common constant of proportionality, the A-parameter. This paradigm is not, however, founded on a compelling underlying theory and therefore it is worthwhile exploring other possibilities.

In this paper we focus on Anomaly Mediation (AM) [1]-[22]. This is a framework in which a single mass parameter determines the $\phi^*\phi$, ϕ^3 and $\lambda\lambda$ supersymmetry-breaking terms in terms of calculable and moreover renormalisation group (RG) invariant functions of the dimensionless couplings, in an elegant and predictive way; too predictive, in fact, in that the theory in its simplest form leads to tachyonic sleptons and fails to accommodate the usual electroweak vacuum state. There is a natural solution to this, however, which restores the correct vacuum while retaining the RG invariance (and hence the ultra-violet insensitivity) of the predictions. This is achieved simply (and without introducing another source of explicit supersymmetry-breaking) by the introduction of a Fayet-Iliopoulos (FI) D-term or terms.

This possibility was first explored in detail in Ref. [17], and subsequently by a number of authors [18]-[22]. The main purpose of this paper is to present the most precise spectrum calculations to date in the AMSB scenario. We also show how the low energy theory employed can arise in a natural way from a theory with an additional anomaly-free U_1 broken at a high scale. We examine the decoupling in this case and show how only the soft mass contributions from the D-terms remain, which can naturally eliminate the tachyonic slepton problem. This provides a minimal extension of anomaly mediation.

In the original scenario of Ref. [17], FI terms corresponding to two distinct U_1 groups were introduced, one being the standard model U_1, and the other the second mixed-anomaly-free (or completely anomaly-free if right-handed neutrinos are included) U_1 admitted by the MSSM. This U_1 may be chosen to be $B-L$ [18], or some linear combination of it and the MSSM U_1 [17]. Or, as emphasised in Ref. [20], a single new U_1 may be employed if the charges are chosen appropriately. If these FI terms are added to the masses with constant coefficients (as in Eq. 4 below) rather than as genuine gauge linear D-terms, then as discussed in Ref. [17] and at more length in Ref [20], the choices made in Refs [17]-[20] are simply reparametrisations of each other. As indicated above, we will see how this scenario can emerge naturally at low energies in a specific theory with an additional gauged anomaly-free U_1.

2 The Minimal Supersymmetric Standard Model

The MSSM is defined by the superpotential:

$$W = H_2 Q Y_t t^c + H_1 Q Y_b b^c + H_1 L Y_\tau \tau^c + \mu H_1 H_2 \tag{1}$$

with soft breaking terms:

$$\begin{aligned} L_{\mathrm{SOFT}} &= \sum_\phi m_\phi^2 \phi^* \phi + \left[m_3^2 H_1 H_2 + \sum_{i=1}^{3} \frac{1}{2} M_i \lambda_i \lambda_i + \mathrm{h.c.} \right] \\ &+ \left[H_2 Q h_t t^c + H_1 Q h_b b^c + H_1 L h_\tau \tau^c + \mathrm{h.c.} \right] \end{aligned} \tag{2}$$

where in general $Y_{t,b,\tau}$ and $h_{t,b,\tau}$ are 3×3 matrices. We work throughout in the approximation that the Yukawa matrices are diagonal, and neglect the Yukawa couplings of the first two generations.

3 The AMSB Solution

Remarkably the following results are RG invariant:

$$\begin{aligned} M_i &= m_0 \beta_{g_i}/g_i \\ h_{t,b,\tau} &= -m_0 \beta_{Y_{t,b,\tau}} \\ (m^2)^i{}_j &= \frac{1}{2} m_0^2 \mu \frac{d}{d\mu} \gamma^i{}_j \\ m_3^2 &= \kappa m_0 \mu - m_0 \beta_\mu \end{aligned} \tag{3}$$

Here β_{g_i} are the gauge β-functions, γ the chiral supermultiplet anomalous dimension, and $\beta_{Y_{t,b,\tau}}$ are the Yukawa β-functions. Moreover, the RG invariance is preserved if we replace $(m^2)^i{}_j$ in Eq. 3 by

$$(\overline{m}^2)^i{}_j = \frac{1}{2} m_0^2 \mu \frac{d}{d\mu} \gamma^i{}_j, + k Y_i \delta^i{}_j, \tag{4}$$

where k is a constant and Y_i are charges corresponding to a U_1 symmetry of the theory with no mixed anomalies with the gauge group. Of course the kY term corresponds in form to a FI D-term. The expressions for M, h and m^2 given in Eq. 3 are obtained if the only source of breaking is a vev in the supergravity multiplet itself: the AMSB scenario (m_0 is then the gravitino mass). Note the parameter κ in the solution for m_3^2; some treatments in the literature omit this term (based on top-down considerations). However, Eq. 3 is RG invariant for arbitrary κ and so we retain it. This means that m_3^2 will be determined in the usual way by the electroweak minimisation. In the following two sections we will show how Eq. 4 can arise via spontaneous breaking of a U_1' symmetry.

Q	u^c	d^c	H_1	H_2	ν^c
$-\frac{1}{3}L$	$-e-\frac{2}{3}L$	$e+\frac{4}{3}L$	$-e-L$	$e+L$	$-2L-e$

Table 1: Anomaly free U_1 symmetry for arbitrary lepton doublet and singlet charges L and e respectively.

4 Anomaly-free U_1 symmetries

The MSSM (including right-handed neutrinos) admits two independent generation-blind anomaly-free U_1 symmetries. The possible charge assignments are shown in Table 1.

Note that an anomaly-free, flavour-blind U_1 necessarily corresponds to equal and opposite charges for $H_{1,2}$ and hence an allowed Higgs μ-term; for an attempt at a Froggatt-Nielsen style origin for the Higgs μ-term using a *flavour dependent* U_1, see for example Ref. [23]. One of the attractive features of anomaly mediation is that squark/slepton mediated flavour changing neutral currents are naturally small; this feature is preserved by a generation-blind U_1 but not by a flavour dependent U_1, so we stick to the former here.

The SM gauged U_1^{SM} is $L=1, e=-2$; this U_1 is of course anomaly free even in the absence of ν^c. U_1^{B-L} is $L=-e=1$; in the absence of ν^c this would have U_1^3 and U_1-gravitational anomalies, but no mixed anomalies with the SM gauge group, which would suffice to maintain the RG invariance of the AMSB solutions. We can introduce FI terms for both U_1^{SM} and U_1^{B-L}, or for U_1^{SM} and a linear combination of them[17], or indeed simply have a single U_1' with the same sign for L and e [20]. We will follow Ref. [20] here; however in the decoupling scenario described in the next section, the low energy theories corresponding to the single U_1' case and and the double U_1 case of Ref. [17] are simply reparametrisations of each other.

5 Spontaneously broken U_1'

With the MSSM augmented by an additional U_1', it is natural to ask at what scale this U_1' is broken. It is possible that this scale is at around 1TeV [24]; here, however, we concentrate on the idea that it is broken at very high energies and that the only low energy remnant of it is the set of FI-type terms that we require.

It would be natural to think that if a U_1' is broken at some high scale M then, by the decoupling theorem, all effects of the U_1' would be suppressed at energies $E << M$ by powers of $1/M$. We shall see that with a FI term this is not the case and it is quite natural for there to be $O(M_{\mathrm{SUSY}})$ scalar mass contributions arising from the presence of the FI term.

It is straightforward to construct a model with an additional gauged U_1' in such a way that the only effect on the low-energy theory is the appearance of the FI terms we require.

We introduce a pair of MSSM singlet fields $\phi, \overline{\phi}$ with U_1' charges $q_{\phi,\overline{\phi}} = \pm(4L+2e)$ and a gauge singlet s, with a superpotential

$$W = \lambda_1 \phi\overline{\phi} s + \frac{1}{2}\lambda_2 \nu^c \nu^c \phi. \tag{5}$$

The choice of charges is essentially determined by the requirement that the $\phi, \overline{\phi}$ fields decouple from the MSSM while generating a large mass for ν^c.

The scalar potential takes the form:

$$\begin{aligned} V &= m_\phi^2 \phi^*\phi + m_{\overline{\phi}}^2 \overline{\phi}^* \overline{\phi} + \cdots \\ &+ \frac{1}{2}\left[\xi - q_\phi(\phi^*\phi - \overline{\phi}^*\overline{\phi}) - \sum_{\text{matter}} e_i \chi_i^* \chi_i\right]^2 + \cdots . \end{aligned} \tag{6}$$

where χ_i stands for all the MSSM scalars, and e_i their U_1' charges, and we have introduced a FI term for U_1'. We will take $\xi > 0$, $q_\phi > 0$ and $\xi >> m_0^2$, and assume that the scalar masses in Eq. 6 (apart from the Higgs μ-term) and other supersymmetry-breaking terms in the theory are the anomaly-mediation contributions. We now proceed to minimise the scalar potential. As we shall see, this will result in a vev for ϕ of order $\sqrt{\xi}$; this means that the appropriate scale at which we should minimise the potential is also around $\sqrt{\xi}$. As a consequence, we of course include at this stage the U_1' contributions in the anomalous dimensions of the fields. It therefore follows that as long as $\lambda_{1,2}$ are somewhat smaller than the U_1' coupling g' then we will have $m_\phi^2 < 0$, which we will assume in the following analysis.

If we look for an extremum with only $\langle\phi\rangle$ nonzero we find

$$\langle\phi^*\phi\rangle \equiv \frac{1}{2}v_\phi^2 = \frac{q_\phi\xi - m_\phi^2}{q_\phi^2} \tag{7}$$

so $\langle\phi\rangle = O(\sqrt{\xi})$ for large ξ and $V \approx m_\phi^2\xi/q_\phi$. Note that since as indicated above we have chosen parameters so that $m_\phi^2 < 0$ we have $V < 0$ at the minimum. Expanding about the minimum, ie with $\phi = (v_\phi + H(x))/\sqrt{2}$, where H is the (real) physical U_1' Higgs, we find

$$\begin{aligned} V &= \frac{m_\phi^2\xi}{q_\phi} - \frac{m_\phi^4}{2q_\phi^2} + (m_\phi^2 + m_{\overline{\phi}}^2 + \frac{1}{2}v_\phi^2\lambda_1^2)\overline{\phi}^*\overline{\phi} - \frac{e_i}{q_\phi}m_\phi^2\chi_i^*\chi_i \\ &+ \frac{1}{2}v_\phi^2\lambda_1^2 s^* s + \frac{1}{2}v_\phi^2\lambda_2^2(\nu^c)^*\nu^c \\ &+ \frac{1}{2}\left(v_\phi q_\phi H - q_\phi\overline{\phi}^*\overline{\phi} + e_i\chi_i^*\chi_i\right)^2 \cdots \end{aligned} \tag{8}$$

For large ξ (i.e. large v_ϕ) all trace of the U_1' in the effective low energy lagrangian disappears, except for contributions to the masses of the matter fields which are naturally of the same order as the AMSB ones. We can see this either by treating the heavy H-field as non-propagating and eliminating it via its equation of motion, or by noting that the quartic $(\chi^*\chi)^2$ D-term still present in Eq. 8 is cancelled (at low energies) by the H-exchange graph using two $H\chi^*\chi$ vertices. In the large ξ limit the breaking of U_1' preserves supersymmetry; thus the U_1' gauge boson, its gaugino, ψ_H and H form a massive supermultiplet which decouples from the theory. The fact that supersymmetry is good at large ξ protects the light χ fields from obtaining masses of $O(\sqrt{\xi})$ from loop corrections. Moreover v_ϕ via the superpotential gives large supersymmetric masses to $\overline{\phi}, s$ and also ν^c , thus naturally implementing the see-saw mechanism. The low energy theory contains just the MSSM fields with the only modification being the FI -type mass contributions proportional to m_ϕ^2 . This is simply another manifestation of the non-decoupling of soft mass corrections from D-terms[25].

We now need only choose the charges L, e for the lepton doublet and singlet so that the contributions to their slepton masses are positive; that is, we choose $L, e > 0$ since $m_\phi^2 < 0$. It is easy to show that (modulo electroweak breaking) this represents the absolute minimum of the potential (note that λ_1 plays a crucial role here in that for $\lambda_1 = 0$ the D-flat direction $\langle\phi\rangle = \langle\overline{\phi}\rangle >> \sqrt{\xi}$ would lead to an potential unbounded from below).

The model constructed here is similar in spirit to those of Harnik et al [19], in that the U_1' breaking is at a high scale so that only the D-term contributions survive in the low energy theory. Like [19] we assume that the anomaly mediated contribution to SUSY breaking is dominant, something that can be justified in the conformal sequestered scheme of Luty et al [16]. The main difference is that we use a FI term to trigger the U_1' breaking rather than an F-term. Here we have just assumed the existence of the FI term as one of the terms allowed by the symmetries of the theory. We will return elsewhere to a discussion of how such a term may be generated in an underlying theory.

In the next section we will explore the region of the (e, L) parameter space such that electroweak-breaking via the Higgses is obtained as usual.

6 The sparticle spectrum

We turn now to the effective low energy theory. Evidently in the scenario described in section 5, we have decoupling of the U_1' at low energies so that the anomalous dimensions of the fields are as in the MSSM; thus for the Higgses and 3rd generation matter fields we have (at one loop):

$$
\begin{aligned}
16\pi^2\gamma_{H_1} &= 3\lambda_b^2 + \lambda_\tau^2 - \tfrac{3}{2}g_2^2 - \tfrac{3}{10}g_1^2, \\
16\pi^2\gamma_{H_2} &= 3\lambda_t^2 - \tfrac{3}{2}g_2^2 - \tfrac{3}{10}g_1^2, \\
16\pi^2\gamma_L &= \lambda_\tau^2 - \tfrac{3}{2}g_2^2 - \tfrac{3}{10}g_1^2, \\
16\pi^2\gamma_Q &= \lambda_b^2 + \lambda_t^2 - \tfrac{8}{3}g_3^2 - \tfrac{3}{2}g_2^2 - \tfrac{1}{30}g_1^2, \\
16\pi^2\gamma_{t^c} &= 2\lambda_t^2 - \tfrac{8}{3}g_3^2 - \tfrac{8}{15}g_1^2, \\
16\pi^2\gamma_{b^c} &= 2\lambda_b^2 - \tfrac{8}{3}g_3^2 - \tfrac{2}{15}g_1^2,
\end{aligned}
$$

$$16\pi^2\gamma_{\tau^c} = 2\lambda_\tau^2 - \tfrac{6}{5}g_1^2, \tag{9}$$

where $\lambda_{t,b,\tau}$ are the third generation Yukawa couplings. For the first two generations we use the same expressions but without the Yukawa contributions.

The soft scalar masses are given by

$$\begin{aligned}
\overline{m}_Q^2 &= m_Q^2 - \tfrac{1}{3}L\xi' \quad \overline{m}_{t^c}^2 = m_{t^c}^2 - (\tfrac{2}{3}L + e)\xi', \\
\overline{m}_{b^c}^2 &= m_{b^c}^2 + (\tfrac{4}{3}L + e)\xi', \quad \overline{m}_L^2 = m_L^2 + L\xi', \\
\overline{m}_{\tau^c}^2 &= m_{\tau^c}^2 + e\xi', \quad \overline{m}_{H_{1,2}}^2 = m_{H_{1,2}}^2 \mp (e + L)\xi',
\end{aligned} \tag{10}$$

where

$$m_Q^2 = \tfrac{1}{2}m_0^2\mu\frac{d}{d\mu}\gamma_Q = \tfrac{1}{2}m_0^2\beta_i\frac{\partial}{\partial\lambda_i}\gamma_Q \tag{11}$$

(where λ_i includes all gauge and Yukawa couplings) and so on, and we have written the effective FI parameter as

$$\xi' = -\frac{m_\phi^2}{q_\phi}. \tag{12}$$

The 3rd generation A-parameters are given by

$$\begin{aligned}
A_t &= -m_0(\gamma_Q + \gamma_{t^c} + \gamma_{H_2}), \\
A_b &= -m_0(\gamma_Q + \gamma_{b^c} + \gamma_{H_1}), \\
A_\tau &= -m_0(\gamma_L + \gamma_{\tau^c} + \gamma_{H_1})
\end{aligned} \tag{13}$$

and we set the corresponding first and second generation quantities to zero. The gaugino masses are given by

$$M_i = m_0|\frac{\beta_{g_i}}{g_i}|. \tag{14}$$

The scale of the FI contributions is set by the AMSB contribution to the ϕ-mass, and hence is naturally expected to be the same order as the other AMSB contributions. In the examples considered below this is indeed the case. Clearly these FI contributions depend on two parameters, $L\xi'$ and $e\xi'$. For notational simplicity we will set $\xi' = 1(\mathrm{TeV})^2$ from now on.

We begin by choosing input values for m_0, $\tan\beta$, L, e and signμ, and then we calculate the appropriate dimensionless coupling input values at the scale M_Z by an iterative procedure involving the sparticle spectrum, and the loop corrections to $\alpha_{1\cdots3}$, m_t, m_b and m_τ, as described in Ref. [27]. We then determine a given sparticle pole mass by running the dimensionless couplings up to a certain scale chosen (by iteration) to be equal to the pole mass itself, and then using Eqs. 11, 13, 14 and including full one-loop corrections from Ref. [27], and two-loop corrections to the top quark mass [26]. As in Ref. [31], we have compared the effect of using one, two and three-loop anomalous dimensions and β-functions in the calculations. Note that when doing the three-loop calculation, we use in Eq. 11, for example, the three loop approximation for both β_i and γ_Q, thus including some higher order effects.

We will present results for $\mu > 0$ and $m_0 = 40\text{TeV}$, for which value the gluino mass is around 900GeV.

The allowed region in (e, L) space corresponding to an acceptable vacuum is shown in Fig. 1. To define the allowed region, we have imposed $m_{\tilde{\tau}} > 82\text{GeV}$, $m_{\tilde{\nu}_\tau} > 49\text{GeV}$ and $m_A > 90\text{GeV}$. The region is to a good approximation triangular, with one side of the triangle corresponding to m_A becoming too light (and quickly imaginary just beyond the boundary, with breakdown of the electroweak vacuum) and the other two sides to one of the sleptons (usually a stau) becoming too light. (Note that Ref. [20] sets $e = 1$ rather than $\xi' = 1$, which is why the allowed region in their Figure 1 has a different shape; the figures are in fact (roughly) equivalent).

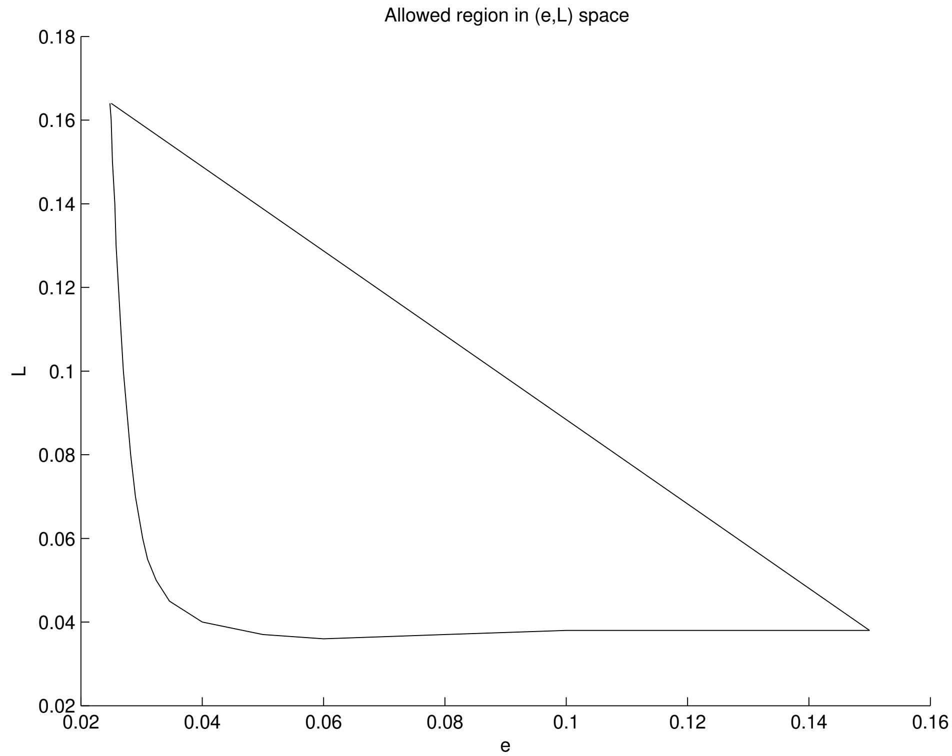

Fig. 1: The region of (e, L) *space corresponding to an acceptable electroweak vacuum, for* $m_0 = 40\,TeV$ *and* $\tan\beta = 10$.

Certain features of the spectrum are apparent from Eq. 10. Since to avoid tachyonic sleptons we must choose $L, e > 0$ we can see that the heaviest squark (especially at low $\tan\beta$) is likely to be the mainly right handed sbottom.

As an example of an acceptable spectrum, we give the results for $m_0 = 40\text{TeV}, \tan\beta = 10, L = 1/25, e = 1/10, \text{sign}\mu = +$ as derived using the one, two and three loop approximations for the anomalous dimensions and β-functions. In Table 2 we have used $m_t = 178\text{GeV}$, while in Table 3 we have used [28] $m_t = 172.7\text{GeV}$. The spectrum is not very much affected by this choice, the most noticeable alteration being in the mass of the light top squark. The rest of the results we present will be for $m_t = 178\text{GeV}$. This point in (e, L) space is near the boundary of the allowed region (see Fig 1) and is characterised by a light stau. The notation in Table 2 etc. is fairly standard, see for example Ref. [27]; note in particular that in all our examples $\tilde{t}_1, \tilde{b}_1, \tilde{\tau}_1$ refer to the mainly lefthanded particles.

The results exhibit the same feature as found for the Snowmass Benchmark (SPS) points in Ref. [31]; that is, the effect of using 3-loop β-functions has a surprisingly large effect on the squark spectrum. This effect was most marked in the SPS case when the gluino mass was significantly larger than the squark masses, which is not the case here; nevertheless, for the light top squark, for example, it is still noticeable.

A characteristic feature of AMSB distinguishing it from MSUGRA is that $M_2 < M_1$, where $M_{1,2}$ are the bino and wino masses respectively. As a result the lightest neutralino (often the LSP) is predominantly the neutral wino and the lighter chargino (often the NLSP) is almost degenerate with it. In Table 2 we have given all results for masses to the nearest GeV; however we have calculated the $\chi^{\pm}$, χ masses using the full one-loop results and expect our results for $\chi_1^{\pm} - \chi_1$ to be good to 10MeV as quoted. A clear account of the dominant contribution to wino mass splitting and the associated phenomenology appears in Ref. [4]. Our splitting of around 240MeV is consistent with their results.

Note that in Table 2 the τ-sneutrino is the LSP; it is interesting that the argument of Ref. [29], which excluded the possibility of a sneutrino LSP in the MSUGRA scenario, does not apply here. The claim was [29] that with MSUGRA boundary conditions, a sneutrino LSP would necessarily have a mass less than half the Z-mass and so contribute to the invisible Z width. Evidently that is not the case here. For a recent discussion of sneutrinos as dark matter see [30].

An interesting feature of the results is that for low $\tan\beta$ we find that the light CP-even Higgs mass, m_h, is less than the experimental (standard model) lower bound of 114GeV. Although the generally quoted supersymmetric bound is significantly lower, we must take seriously the SM bound here, since we find generally that for us $\sin(\beta-\alpha) \sim 1$, so that h couples to the Z-boson like the SM Higgs. However as $\tan\beta$ is increased, m_h increases above this bound. The allowed range of $\tan\beta$ depends on the choice of L, e. In Fig 2,3 we plot m_h and the CP-odd Higgs mass m_A against $\tan\beta$ for $L = 1/25, e = 1/10$. The electroweak vacuum fails for $\tan\beta > 25$ in this case. We also plot the lighter stau mass (Fig. 4) and the tau sneutrino mass (Fig. 5) against $\tan\beta$. We see that acceptable values of m_h are obtained for $7 < \tan\beta < 25$, and of the stau mass for $\tan\beta < 19$. As can be seen from Table 2, m_h is actually essentially unchanged by whether we use two or three-loop β-functions; in fact we have used the two-loop β-functions to generate Figs. 2,3.

In Table 4, we give the results for another point in (e, L)-space, chosen to be in the centre of the allowed region, where this time the lightest neutralino is the LSP.

Finally in Table 5 we give results for $(e, L) = (0.05, 0.05)$, a point again near the boundary in (e, L) space, with light sleptons and heavy squarks, and also a large charged Higgs mass of over 400GeV. This point is interesting because of the fact that previous authors have noted that the fact that M_3 and M_2 have opposite signs disfavours at first sight a supersymmetric explanation of the well-known discrepancy between theory and experiment for the anomalous magnetic moment of the muon,

mass (GeV)	1loop	2loops	3loops
$\tilde{g}$	914	891	888
$\tilde{t}_1$	770	761	751
$\tilde{t}_2$	543	540	529
$\tilde{u}_L$	834	820	809
$\tilde{u}_R$	767	756	744
$\tilde{b}_1$	738	728	718
$\tilde{b}_2$	928	920	910
$\tilde{d}_L$	838	824	813
$\tilde{d}_R$	937	929	919
$\tilde{\tau}_1$	124	109	110
$\tilde{\tau}_2$	284	284	284
$\tilde{e}_L$	132	118	119
$\tilde{e}_R$	285	285	285
$\tilde{\nu}_e$	104	86	86
$\tilde{\nu}_\tau$	99	79	80
χ_1	105	129	129
χ_2	354	362	361
χ_3	540	563	555
χ_4	552	575	566
$\chi_1^\pm$	106	129	129
$\chi_2^\pm$	549	572	563
h	117	117	117
H	315	351	336
A	314	351	336
$H^\pm$	324	360	345
$\chi_1^\pm - \chi_1$ (MeV)	230	240	240

Table 2: Mass spectrum for $m_t = 178\text{GeV}$, $m_0 = 40\text{TeV}$, $\tan\beta = 10$, $L = 1/25$, $e = 1/10$

mass (GeV)	1loop	2loops	3loops
$\tilde{g}$	914	890	888
$\tilde{t}_1$	767	758	748
$\tilde{t}_2$	519	516	505
$\tilde{u}_L$	835	820	809
$\tilde{u}_R$	767	756	744
$\tilde{b}_1$	732	723	713
$\tilde{b}_2$	928	920	910
$\tilde{d}_L$	838	824	813
$\tilde{d}_R$	937	929	919
$\tilde{\tau}_1$	123	108	109
$\tilde{\tau}_2$	284	284	284
$\tilde{e}_L$	132	118	119
$\tilde{e}_R$	285	285	285
$\tilde{\nu}_e$	104	85	86
$\tilde{\nu}_\tau$	99	79	80
χ_1	106	129	130
χ_2	355	362	362
χ_3	563	584	576
χ_4	574	595	587
$\chi_1^\pm$	106	130	130
$\chi_2^\pm$	571	592	584
h	116	115	115
H	354	385	372
A	353	384	372
$H^\pm$	362	393	381
$\chi_1^\pm - \chi_1$ (MeV)	220	230	230

Table 3: Mass spectrum for $m_t = 172.7\text{GeV}$, $m_0 = 40\text{TeV}$, $\tan\beta = 10$, $L = 1/25$, $e = 1/10$

mass (GeV)	1loop	2loops	3loops
$\tilde{g}$	914	890	888
$\tilde{t}_1$	762	753	743
$\tilde{t}_2$	554	541	530
$\tilde{u}_L$	826	812	801
$\tilde{u}_R$	769	758	746
$\tilde{b}_1$	728	719	709
$\tilde{b}_2$	940	932	923
$\tilde{d}_L$	830	816	805
$\tilde{d}_R$	949	941	932
$\tilde{\tau}_1$	212	208	208
$\tilde{\tau}_2$	250	247	247
$\tilde{e}_L$	228	228	228
$\tilde{e}_R$	241	234	235
$\tilde{\nu}_e$	227	220	220
$\tilde{\nu}_\tau$	225	218	218
χ_1	106	130	130
χ_2	353	361	361
χ_3	530	554	545
χ_4	543	566	557
$\chi_1^\pm$	106	130	130
$\chi_2^\pm$	539	562	553
h	117	117	117
H	277	319	303
A	277	319	302
$H^\pm$	288	329	313
$\chi_1^\pm - \chi_1$ (MeV)	240	250	250

Table 4: Mass spectrum for $m_0 = 40$TeV, $\tan\beta = 10$, $L = 0.08$, $e = 0.07$

a_μ. This is because if sign (μM_2) is chosen so as to create a positive $a_\mu^{\rm SUSY}$ then sign (μM_3) leads to constructive interference between various supersymmetric contributions to $B(b \to s\gamma)$, and consequent restrictions on the allowed parameter space. However, with light sleptons (to generate a contribution to a_μ) and heavy squarks and charged Higgs (to suppress contributions to $B(b \to s\gamma)$) this conclusion can be evaded (as was already argued in Ref. [32]).

There has been a considerable amount of work in recent years on two loop corrections to m_h [33]-[36]. For some regions of the MSSM parameter space these can be substantial; therefore since we have presented predictions of around $115 - 118$GeV we have to worry about them since they generally reduce m_h. Using the useful web resource from Ref. [33], we obtain, for the input parameters of Table 5, the result $m_h = 116.2 \pm 1.4$GeV, in excellent agreement with our results, which suggests that the two-loop corrections are in fact not very large in our scenario. Other points in (e, L) space give similar results. Thus for $m_0 = 40$TeV we predict that m_h is less than about 118.4GeV (see Fig. 2). If we increase m_0 then this bound does increase somewhat (to around 125GeV at $m_0 = 100$TeV, for example) but at the price of considerable electroweak fine-tuning.

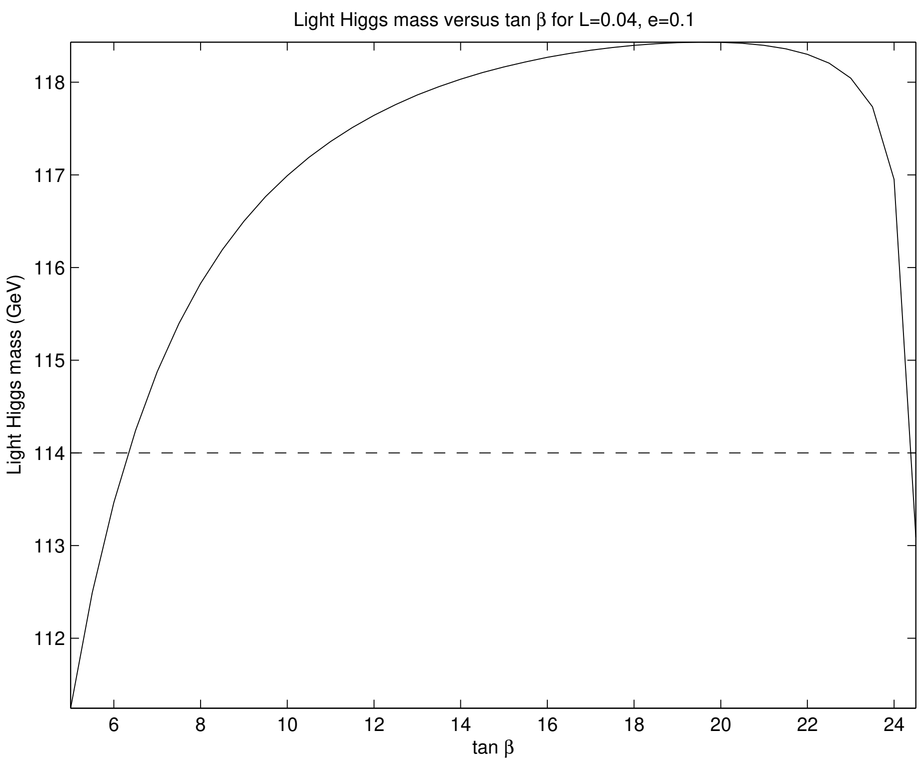

Fig. 2: The light CP-even Higgs mass m_h as a function of $\tan\beta$, for $L = 1/25, e = 1/10$. The dotted line is the SM lower limit (114GeV).

mass (GeV)	1loop	2loops	3loops
$\tilde{g}$	914	890	888
$\tilde{t}_1$	772	763	753
$\tilde{t}_2$	579	576	566
$\tilde{u}_L$	832	818	807
$\tilde{u}_R$	795	785	773
$\tilde{b}_1$	737	727	718
$\tilde{b}_2$	909	900	890
$\tilde{d}_L$	836	822	811
$\tilde{d}_R$	917	909	899
$\tilde{\tau}_1$	140	130	131
$\tilde{\tau}_2$	194	191	191
$\tilde{e}_L$	168	158	158
$\tilde{e}_R$	178	177	177
$\tilde{\nu}_e$	148	137	137
$\tilde{\nu}_\tau$	144	133	133
χ_1	106	130	130
χ_2	355	362	362
χ_3	577	599	590
χ_4	587	609	601
$\chi_1^\pm$	106	130	130
$\chi_2^\pm$	585	606	598
h	117	117	117
H	429	455	444
A	428	454	443
$H^\pm$	436	462	451
$\chi_1^\pm - \chi_1$ (MeV)	220	230	230

Table 5: Mass spectrum for $m_0 = 40\text{TeV}$, $\tan\beta = 10$, $L = 0.05$, $e = 0.05$

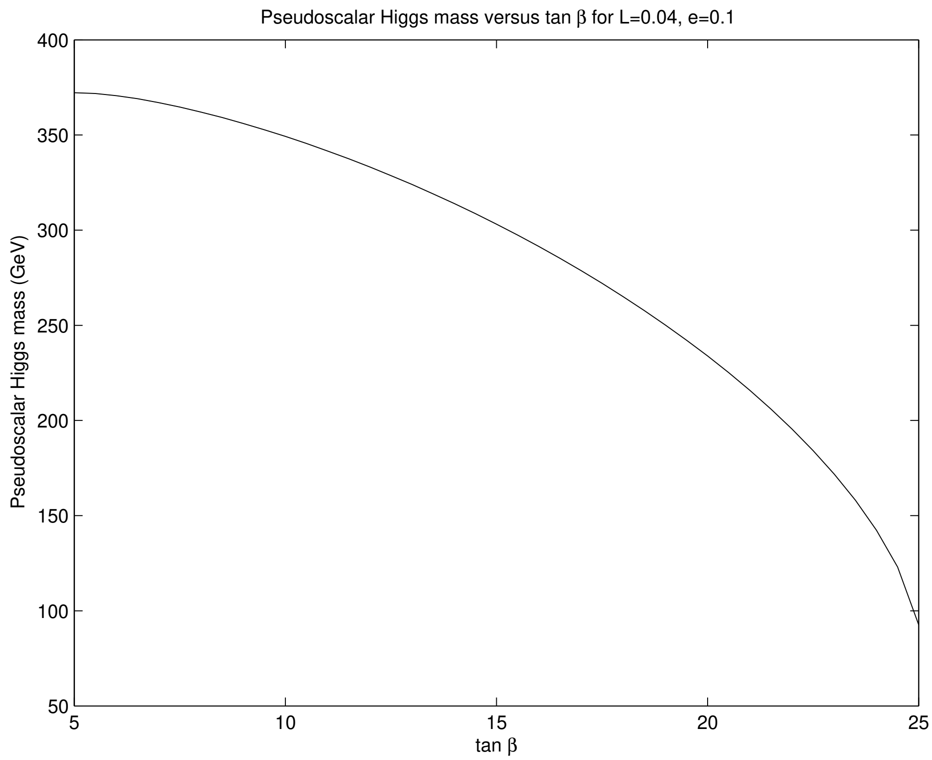

Fig. 3: The CP-odd Higgs mass m_A as a function of $\tan\beta$, for $L = 1/25, e = 1/10$.

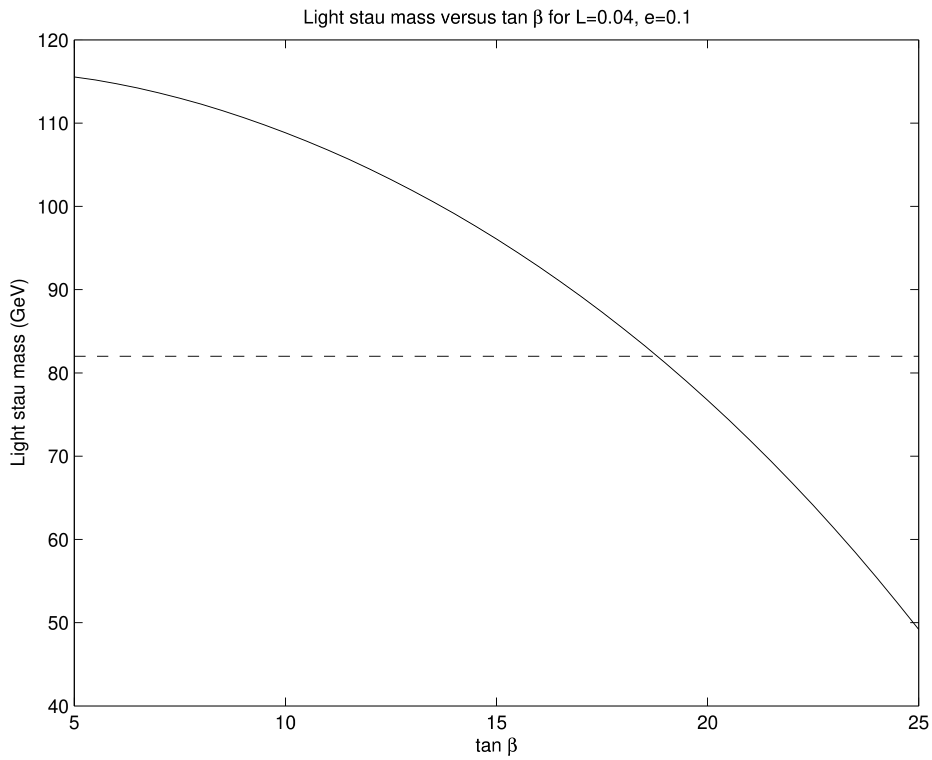

Fig. 4: The light stau mass as a function of $\tan\beta$*, for* $L = 1/25, e = 1/10$*. The dotted line is the lower limit* $(82GeV)$*.*

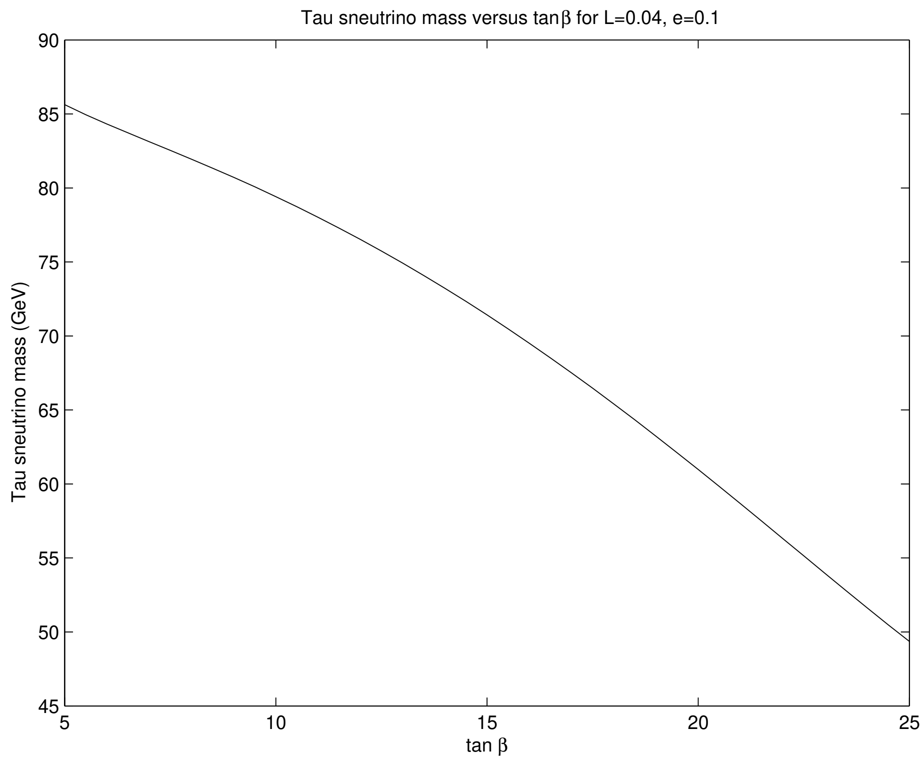

Fig. 5: The τ-sneutrino mass as a function of $\tan\beta$*, for* $L = 1/25, e = 1/10$.

7 Mass sum rules

By taking appropriate linear combinations of masses it is straightforward to derive a set of interesting sum rules [17].

In the following equations, if we substitute the tree values for the various masses on the left hand side, the (e, L) dependent terms and the electroweak breaking contributions to the masses cancel. We have calculated the numerical coefficients on the right hand side from Table 2, using the two-loop sparticle mass predictions; it is easy to then check that to the indicated accuracy the same equations hold for the results in Tables 4, 5. Thus these sum rules are to an excellent approximation independent of (e, L), and also in fact of m_0; the numerical coefficients are slowly varying functions of $\tan\beta$ and the input top pole mass.

$$\begin{aligned} m^2_{\tilde{t}_1} + m^2_{\tilde{t}_2} + m^2_{\tilde{b}_1} + m^2_{\tilde{b}_2} - 2m_t^2 &= 2.76\,(m_{\tilde{g}})^2 \\ m^2_{\tilde{\tau}_1} + m^2_{\tilde{\tau}_2} + m^2_{\tilde{t}_1} + m^2_{\tilde{t}_2} - 2m_t^2 &= 1.14\,(m_{\tilde{g}})^2 . \end{aligned} \tag{15}$$

$$m^2_{\tilde{e}_L} + 2m^2_{\tilde{u}_L} + m^2_{\tilde{d}_L} = 2.60\,(m_{\tilde{g}})^2 ,$$

$$
\begin{aligned}
m^2_{\tilde{u}_R} + m^2_{\tilde{d}_R} + m^2_{\tilde{u}_L} + m^2_{\tilde{d}_L} &= 3.51\,(m_{\tilde{g}})^2\,, \\
m^2_{\tilde{u}_L} + m^2_{\tilde{d}_L} - m^2_{\tilde{u}_R} - m^2_{\tilde{e}_R} &= 0.88\,(m_{\tilde{g}})^2\,.
\end{aligned} \tag{16}
$$

$$
\begin{aligned}
m^2_A - 2\sec 2\beta \left(m^2_{\tilde{e}_L} + m^2_{\tilde{e}_R}\right) &= 0.40\,(m_{\tilde{g}})^2\,, \\
m^2_A - 2\sec 2\beta \left(m^2_{\tilde{\tau}_1} + m^2_{\tilde{\tau}_2} - 2m^2_{\tau}\right) &= 0.39\,(m_{\tilde{g}})^2\,.
\end{aligned} \tag{17}
$$

The existence of these sum rules will be a useful distinguishing feature of the AMSB scenario.

8 Conclusions

Despite remarkable advances in the understanding of string theory, a coherent high energy theory spawning the MSSM as an effective low energy theory remains elusive. This has led to exploration of such outré possibilities as little higgs models and split supersymmetry. Remaining within the conservative world of low energy supersymmetry, the AMSB scenario is an attractive alternative to (and easily distinguished from) MSUGRA. We have shown how a U_1' gauge symmetry broken at high energies can lead in a natural way to the FI-solution to the tachyonic slepton problem in the context of anomaly mediation. The result is a sparticle spectrum described by the parameter set $m_0, e, L, \tan\beta, \mathrm{sign}(\mu)$; and it is only for a comparatively restricted set of (e, L) that an acceptable spectrum is obtained. Moreover we have presented a set of sum rules which are independent of m_0, L and e. At the very least, the scenario we describe has the merit of being immediately testable should sparticles be discovered in experiments at the LHC.

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References

[1] L. Randall and R. Sundrum, *Nucl. Phys.* B557 (1999) 79

[2] G.F. Giudice et al, *JHEP* 9812 (1998) 27

[3] A. Pomarol and R. Rattazzi, *JHEP* 9905 (1999) 013

[4] T. Gherghetta, G.F. Giudice and J.D. Wells, *Nucl. Phys.* B559 (1999) 27

[5] M.A. Luty and R. Rattazzi, *JHEP* 9911 (1999) 001

[6] Z. Chacko, M.A. Luty, I. Maksymyk and E. Ponton, *JHEP* 0004 (2000) 001

[7] E. Katz, Y. Shadmi and Y. Shirman, *JHEP* 9908 (1999) 015

[8] I. Jack and D.R.T. Jones, *Phys. Lett.* B465 (1999) 148

[9] J.L. Feng and T. Moroi, *Phys. Rev.* D61 (2000) 095004

[10] G.D. Kribs, *Phys. Rev.* D62 (2000) 015008

[11] S. Su, *Nucl. Phys.* B573 (2000) 87

[12] J.A. Bagger, T. Moroi and E. Poppitz, *JHEP* 0004 (2000) 009

[13] R. Rattazzi et al, *Nucl. Phys.* B576 (2000) 3

[14] F.E. Paige and J. Wells, hep-ph/0001249

[15] N. Okada, *Phys. Rev.* D65 (2002) 115009

[16] M. Luty and R. Sundrum, *Phys. Rev.* D67 (2003) 045007

[17] I. Jack and D.R.T. Jones, *Phys. Lett.* B482 (2000) 167

[18] N. Arkani-Hamed, D. E. Kaplan, H. Murayama and Y. Nomura, *JHEP* 0102 (2001) 041

[19] R. Harnik, H. Murayama and A. Pierce, *JHEP* 0208 (2002) 034

[20] B. Murakami and J. D. Wells, *Phys. Rev.* D68 (2003) 035006

[21] R. Kitano, G. D. Kribs and H. Murayama, *Phys. Rev.* D70 (2004) 035001

[22] M. Ibe, R. Kitano and H. Murayama, *Phys. Rev.* D71 (2005) 075003

[23] I. Jack, D.R.T. Jones and R. Wild, *Phys. Lett.* B580 (2004) 72

[24] J. Erler and P. Langacker, *Phys. Rev, Lett.* 84 (2000) 212;
D.A. Demir, G.L. Kane and T.T. Wang, hep-ph/0503290.

[25] Y. Kawamura, H. Murayama and M. Yamaguchi, *Phys. Rev.* D51 (1995) 1337;
H. Murayama, hep-ph/9503392

[26] A. Bednyakov et al, *Eur. Phys. J.* C29 (2003) 87

[27] D.M. Pierce, J.A. Bagger, K.T. Matchev and R.J. Zhang, *Nucl. Phys.* B491 (1997) 3

[28] The CDF Collaboration, the D0 Collaboration and the TEVEWWWG, hep-ex/0507091

[29] T. Hebbeker, *Phys. Lett.* B470 (1999) 259

[30] D. Hooper, J. March-Russell and S.M. West, *Phys. Lett.* B605 (2005) 228

[31] I. Jack, D.R.T. Jones, A.F. Kord, *Ann. Phys.* 316 (2005) 213

[32] I. Jack and D.R.T. Jones *Nucl. Phys.* B662 (2003) 63

[33] S. Heinemeyer, W. Hollik and G. Weiglein, *Eur. Phys. Jour.* C 9 (1999) 343, *Comp. Phys. Comm.* 124 2000 76;
The code is accessible via `www.feynhiggs.de`

[34] J. R. Espinosa and R. J. Zhang, *Nucl. Phys.* B586 (2000) 3; *JHEP* 0003 2000 026

[35] A. Brignole et al, *Nucl. Phys.* B631 (2002) 195; *ibid* B643 (2002) 79

[36] S. P. Martin, *Phys. Rev.* D67 (2003) 095012; *ibid* 67 (2003) 095012