

Fayet-Iliopoulos D -terms, neutrino masses and anomaly mediated supersymmetry breaking

I. Jack, D.R.T. Jones and R. Wild

Dept. of Mathematical Sciences, University of Liverpool, Liverpool L69 3BX, U.K.

We explore, in the context of the MSSM generalised to admit massive neutrinos, an extension of the Anomaly Mediated Supersymmetry Breaking solution for the soft scalar masses to incorporate Fayet-Iliopoulos D -terms. The slepton mass problem characteristic of the scenario is resolved, and the fermion mass hierarchy is explained via the Froggatt-Nielsen mechanism. FCNC problems are evaded by a combination of universal doublet charges and Yukawa textures which are diagonalised by transforming the left-handed fields only.

Recently there has been interest in a specific and predictive framework for the origin of soft supersymmetry breaking within the MSSM, known as Anomaly Mediated Supersymmetry Breaking (AMSB) [1]–[34]. Direct application of the AMSB solution to the MSSM leads, unfortunately, to negative (mass)² sleptons. A number of possible solutions to this problem have been discussed; here we concentrate on a proposal by two of us [13] (see also [26]); namely the introduction of Fayet-Iliopoulos (FI) terms associated with both the MSSM U_1 and an additional U'_1 symmetry. This preserves the exact RG invariance of the AMSB solution in rather a minimalist way, requiring as it does the introduction of no new fields; the U'_1 need not in fact be gauged, though the RG invariance requires that we ensure that it has vanishing linear mixed anomalies with the MSSM gauge group. The MSSM indeed admits two generation-independent, mixed-anomaly-free U_1 groups, the existing U_1^Y and another (which could be chosen to be U_1^{B-L} [26], or a linear combination of it and U_1^Y).

Extension of this scenario to include massive neutrinos meets an obstacle inasmuch as there is no flavour independent global U'_1 symmetry possible for a superpotential incorporating both neutrino Yukawa couplings and Majorana masses for right-handed neutrinos. We are therefore driven to consider flavour dependent symmetries¹, and choose to make a virtue out of necessity in this regard by re-examining the well-travelled path of Yukawa coupling textures associated with a U_1 symmetry [35]. Most of the (considerable) literature on this subject has dealt with an *anomalous* U'_1 (with anomaly cancellation finally achieved via the Green-Schwarz mechanism [36]). This route is not open to us, however, as we require cancellation of the mixed gauge- U'_1 anomalies. We believe, however, that the conclusion that texture generation via an anomaly-free U'_1 is impossible is a consequence of assumptions which, while plausible, are not strictly necessary (what constitutes necessity being in these matters, we admit, a question of taste).

Thus our goal is to show that the MSSM with massive neutrinos (which we will call the MSSM ^{ν}) admits a global U'_1 which enables us to solve the AMSB tachyonic slepton problem while simultaneously reproducing the well-known hierarchies [37]

$$m_\tau : m_\mu : m_e = m_b : m_s : m_d = 1 : \lambda^2 : \lambda^4, \quad \text{and} \quad m_t : m_c : m_u = 1 : \lambda^4 : \lambda^8, \quad (1)$$

where $\lambda \approx 0.22$.

¹ for some alternative ideas see Ref. [26]

We will assume that Yukawa terms are generated via the Froggatt-Nielsen (FN) mechanism[35]: specifically, from higher dimension terms involving MSSM $^\nu$ singlet fields $\theta_{t,b,\tau}$ with each θ associated with a *particular* Yukawa matrix, via terms such as $H_2 Q_i t_j^c (\frac{\theta_t}{M_U})^{a_{ij}}$, where M_U represents the scale of new physics. Then if we require Yukawa textures consistent with the above hierarchies, and also require that the mixed anomalies cancel, we are in general led to consider different charges for each of $\theta_{t,b,\tau}$. If we choose θ -charges

$$q_{\theta_t} = -1, \quad q_{\theta_b} = 2 - \frac{\Delta}{2}, \quad q_{\theta_\tau} = 2 - \frac{\Delta}{3}, \quad (2)$$

where $\Delta = h_1 + h_2$, and the charge assignments shown in Table 1,

Q_i	t_2^c	t_3^c	b_1^c	b_2^c	b_3^c
$8 - t_1^c - h_2$	$t_1^c - 4$	$t_1^c - 8$	$3h_2 + h_1 + t_1^c - 16$	$2h_2 + t_1^c - 12$	$t_1^c - h_1 + h_2 - 8$

L_i	τ_1^c	τ_2^c	τ_3^c
$3t_1^c - \frac{1}{3}h_1 + \frac{8}{3}h_2 - 24$	$\frac{2}{3}h_1 - \frac{4}{3}h_2 + 16 - 3t_1^c$	$20 - 2h_2 - 3t_1^c$	$24 - \frac{2}{3}h_1 - \frac{8}{3}h_2 - 3t_1^c$

Table 1: The U'_1 -charges

then the mixed anomalies cancel and we find textures given by

$$Y_t \sim \begin{pmatrix} \lambda^8 & \lambda^4 & 1 \\ \lambda^8 & \lambda^4 & 1 \\ \lambda^8 & \lambda^4 & 1 \end{pmatrix}, \quad Y_b \sim \begin{pmatrix} \lambda^4 & \lambda^2 & 1 \\ \lambda^4 & \lambda^2 & 1 \\ \lambda^4 & \lambda^2 & 1 \end{pmatrix}, \quad Y_\tau \sim \begin{pmatrix} \lambda^4 & \lambda^2 & 1 \\ \lambda^4 & \lambda^2 & 1 \\ \lambda^4 & \lambda^2 & 1 \end{pmatrix}, \quad (3)$$

where we assume $\langle \theta_{t,b,\tau} \rangle / M_U \approx \lambda \approx 0.22$. The powers of λ are determined by relations such as $h_1 + Q_i + b_1^c + 4q_{\theta_b} = 0$. The equality of the rows in each matrix corresponds to generation independent doublet charges Q_i and L_i . (We use the same notation for the field and its U'_1 charge; it should be clear from the context which is intended.)

Textures of this form have in fact been considered before in the context of D-branes [38], and termed ‘‘single right-handed democracy’’. It is easy to show that the eigenvalues of $Y_{t,b,\tau}$ above lead to the mass textures of Eq. (1). Another feature of textures of this generic form is that since to a good approximation we have

$$Y_t^T Y_t \sim \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix} \quad (4)$$

(similarly for $Y_{b,\tau}$) it is evident that the rotation on the RH fields required to diagonalise the mass matrix will be of the generic form

$$U_R \sim \begin{pmatrix} \cos \alpha & \sin \alpha & 0 \\ -\sin \alpha & \cos \alpha & 0 \\ 0 & 0 & 1 \end{pmatrix}. \quad (5)$$

Moreover, by considering the specific textures shown in Eq. (3) in the approximation that we set to zero the first column of each matrix it is easy to show that in our case U_R will always be close to the unit matrix[38]. This will be significant later when we consider flavour changing neutral currents (FCNCs). If we assume the specific forms

$$Y_t \propto \begin{pmatrix} a_t \lambda^8 & d_t \lambda^4 & 1 + O(\lambda^2) \\ b_t \lambda^8 & e_t \lambda^4 & 1 + O(\lambda^2) \\ c_t \lambda^8 & f_t \lambda^4 & 1 + O(\lambda^2) \end{pmatrix} \quad \text{and} \quad Y_b \propto \begin{pmatrix} a_b \lambda^4 & d_b \lambda^2 & 1 + O(\lambda^2) \\ b_b \lambda^4 & e_b \lambda^2 & 1 + O(\lambda^2) \\ c_b \lambda^4 & f_b \lambda^2 & 1 + O(\lambda^2) \end{pmatrix} \quad (6)$$

then we obtain for the CKM matrix the texture

$$CKM \sim \begin{pmatrix} 1 & 1 & \lambda^2 \\ 1 & 1 & \lambda^2 \\ \lambda^2 & \lambda^2 & 1 \end{pmatrix} \quad (7)$$

which is not of the form of the standard Wolfenstein parametrisation,

$$CKM_W \sim \begin{pmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix} \quad (8)$$

It does, however, reproduce the most significant feature, which is the smallness of the couplings to the third generation.² It follows that it is possible to exhibit explicit forms of Y_t, Y_b with “reasonable” coefficients reproducing the observed CKM matrix, for example:

$$Y_t \propto \begin{pmatrix} 11.35\lambda^8 & 0.915\lambda^4 & 1.048 \\ 1.244\lambda^8 & 3.336\lambda^4 & 0.970 \\ 3.362\lambda^8 & -4.266\lambda^4 & 0.980 \end{pmatrix} \quad Y_b \propto \begin{pmatrix} 0.487\lambda^4 & 0.281\lambda^2 & 1.063 \\ -1.311\lambda^4 & 0.398\lambda^2 & 1.008 \\ -0.514\lambda^4 & -0.750\lambda^2 & 0.925 \end{pmatrix}. \quad (9)$$

The ub and td entries in the CKM matrix are comparatively sensitive to changes in the coefficients in Eq. (9), because our texture form, Eq. (6), does not *naturally* explain the factor of 10 difference between these entries and the cb and ts ones.

² We disagree somewhat with Ref. [38], where it is asserted that the Wolfenstein texture follows if we replace the (13) elements of both Y_t and Y_b in Eq. (6) by $1 + O(\lambda)$. With this particular form for $Y_{t,b}$ it is straightforward to establish (by either numerical or analytic means) that the CKM matrix would have a texture similar to Eq. (7) but with λ^2 replaced everywhere by λ .

The charges shown in Table 1 have been chosen to provide cancellation of the mixed $U'_1(SU_3)^2$, $U'_1(SU_2)^2$, and $U'_1(U_1)^2$ anomalies. This is what is required to render our scalar mass solution RG invariant; but it is also of interest to examine the remaining anomalies involving U'_1 . The $(U'_1)^2 U_1$, $(U'_1)^3$ and (U'_1) -gravitational anomalies are proportional respectively to

$$A_Q = \Delta \left[8h_2 + 6t_1^c - \frac{224}{3} + \frac{14}{9} \Delta \right], \quad (10)$$

$$\begin{aligned} A_C = & -108t_1^c \Delta h_2 - 2368\Delta - \frac{632}{3} \Delta^2 + 640\Delta h_2 + 816t_1^c \Delta + \frac{92}{3} h_2 \Delta^2 - 48\Delta h_2^2 \\ & + 16t_1^c \Delta^2 - 54(t_1^c)^2 \Delta + \frac{16}{9} \Delta^3 \\ & + 3(9(t_1^c)^2 - 168t_1^c + 24h_2 t_1^c + 16h_2^2 + 816 - 224h_2)(3t_1^c + 4h_2 - 28), \end{aligned} \quad (11)$$

and

$$A_G = 3(3t_1^c + 4h_2 - 28). \quad (12)$$

Note that there will be additional contributions to A_G and A_C from any MSSM' singlet fields with U'_1 charges, such as, of course, the θ -fields introduced above, unless they are accompanied by oppositely charged partners. In the specific case of the θ -fields, such $\bar{\theta}$ -partners (if they exist) must be forbidden from generating textures in order to preserve the patterns of Eq. (3). One reason for assuming the $\bar{\theta}$ s exist is that unless they do (and have vevs approximately equal to the corresponding θ s) then the quadratic D -terms for the U'_1 (if it is gauged) will generate large masses for all the MSSM' fields [39]. The possible generation of weak-scale contributions to the sparticle masses via this route and their impact on flavour-changing processes was discussed recently in Ref. [40].

For our purpose it is not necessary to gauge U'_1 , or remove its anomalies other than the linear mixed ones; let us, however, explore (en passant) the option of rendering it completely anomaly-free. It is easy to show that the conditions

$$\Delta = 0 \quad \text{and} \quad 3t_1^c + 4h_2 - 28 = 0 \quad (13)$$

are necessary and sufficient to give $A_Q = A_C = A_G = 0$ above. (Cancellation of all the anomalies requires the existence of the θ -partners $\bar{\theta}$ as described above; if the θ -contributions in A_C and A_G are not cancelled, it is straightforward to show there is no charge assignment such that $A_Q = A_C = A_G = 0$.) The solution Eq. (13) may be of interest in the non-AMSB context, providing as it does a fully anomaly free U'_1 (and one

which would remain anomaly-free with the introduction of $SU_3 \otimes SU_2 \otimes U_1$ singlets in \pm pairs). However it turns out that we cannot with these conditions resolve the slepton mass problem, as we shall now show.

Including the FI contributions the scalar masses are given by[13]

$$\begin{aligned}
\overline{m}_Q^2 &= m_Q^2 + \frac{1}{6}\zeta_1 + \zeta_2 Q_i \delta_j^i, & \overline{m}_{t^c}^2 &= m_{t^c}^2 - \frac{2}{3}\zeta_1 + \zeta_2 t_i^c \delta_j^i, \\
\overline{m}_{b^c}^2 &= m_{b^c}^2 + \frac{1}{3}\zeta_1 + \zeta_2 b_i^c \delta_j^i, & \overline{m}_L^2 &= m_L^2 - \frac{1}{2}\zeta_1 + \zeta_2 L_i \delta_j^i, \\
\overline{m}_{\tau^c}^2 &= m_{\tau^c}^2 + \zeta_1 + \zeta_2 \tau_i^c \delta_j^i, & \overline{m}_{H_1}^2 &= m_{H_1}^2 - \frac{1}{2}\zeta_1 + \zeta_2 h_1, \\
\overline{m}_{H_2}^2 &= m_{H_2}^2 + \frac{1}{2}\zeta_1 + \zeta_2 h_2
\end{aligned} \tag{14}$$

where $\zeta_{1,2}$ are constants and where m_Q^2, \dots denote the AMSB contributions. For example[1]-[3],

$$(m_Q^2)^i_j = \frac{1}{2}|m_0|^2 \mu \frac{d(\gamma_Q)^i_j}{d\mu}, \tag{15}$$

where γ_Q is the quark doublet anomalous dimension matrix and m_0 is the gravitino mass. The slepton mass problem is the fact that m_L^2 and $m_{\tau^c}^2$ have negative eigenvalues. However, as we shall see, we can choose U_1' charges so that for some region of $\zeta_{1,2}$ parameter space, the eigenvalues of \overline{m}_L^2 and $\overline{m}_{\tau^c}^2$ are all positive, and indeed we obtain a fully realistic mass spectrum.

In fact, it is easily shown that with the charge assignments shown in Table 1, and $\Delta = 0$, there exists some range of $\zeta_{1,2}$ leading to positive FI contributions for both $\overline{m}_{\tau^c}^2$ and \overline{m}_L^2 if and only if

$$3t_1^c + 4h_2 < 24 \quad \text{or} \quad 3t_1^c + 4h_2 > 32, \tag{16}$$

so that the fully anomaly-free solution Eq. (13) is excluded. We will nevertheless choose to impose cancellation of the quadratic (A_Q) anomaly since (unlike A_C and A_G) it cannot be affected by a $SU_3 \otimes SU_2 \otimes U_1$ singlet sector. Obviously for $A_Q = 0$ we require either $\Delta = 0$ or $8h_2 + 6t_1^c - \frac{224}{3} + \frac{14}{9}\Delta = 0$. We start with the $\Delta = 0$ case, postponing discussion of the second possibility until later. On the one hand this means that a Higgs μ -term is allowed, and so we have no solution to the μ -problem; on the other hand from Eq. (2) we see that the same θ -field will in fact serve for both down and charged lepton matrices. Moreover we do not need to forbid terms of the kind, for example, $H_2 Q_i t_j^c (\frac{\theta_b}{M_U})^{a_{ij}}$, since no such (gauge invariant) term can be constructed.

We turn now to the issue of neutrino masses. If we wish to generate them via the seesaw mechanism then this suggests that we introduce three right-handed neutrinos with

zero U'_1 charges; however to obtain Dirac mass terms of the form $H_2 L \nu^c$ we then require a θ -field capable of being matched to $L_i + h_2 = 3t_1^c + 4h_2 - 24$. But if we examine potential dimension 4 R -parity violating operators, we find that the possible superpotential operators $t_2^c b_2^c b_3^c$, $t_1^c b_1^c b_3^c$, $Q_i L_j b_3^c$ and $L_i L_j \tau_3^c$ have the same charge as $L_i h_2$, with disastrous consequences for proton decay, if we allow them to be generated at the same level as the $H_2 L \nu^c$ terms. We could choose to impose R -parity conservation, but an attractive alternative is to introduce only two right-handed neutrinos, with charges $\nu_{1,2}^c = \pm q_\nu$ and introduce θ_ν with charge q_{θ_ν} such that

$$L_i + h_2 + q_\nu + n q_{\theta_\nu} = 0, \quad \text{and} \quad L_i + h_2 - q_\nu + m q_{\theta_\nu} = 0 \quad (17)$$

for integer m, n . It is easy to show that if we choose, for example, $n = 2$ and $m = 1$ and $q_{\theta_\nu} = -9$, then none of the R -parity violating Yukawas mentioned above (nor any R -violating Yukawa) can be generated using the available θ -charges. Consequently unrealistic proton decay is prevented. We obtain a ν^c matrix of the form (for consideration of various forms for the ν^c matrix see for example [41])

$$M_{\nu^c} = \begin{pmatrix} 0 & M_1^\nu \\ M_1^\nu & 0 \end{pmatrix} \quad (18)$$

which has non-zero determinant and therefore will serve for the seesaw.³ Moreover the Dirac matrix from $H_2 L \nu^c$ takes the form

$$m_D = \begin{pmatrix} a_\nu \lambda^n & d_\nu \lambda^m \\ b_\nu \lambda^n & e_\nu \lambda^m \\ c_\nu \lambda^n & f_\nu \lambda^m \end{pmatrix} v_2, \quad (19)$$

(where $\langle H_2 \rangle = \begin{pmatrix} 0 \\ v_2 \end{pmatrix}$) and the eigenvalues of the resulting LH neutrino mass matrix

$$m_\nu = m_D M_{\nu^c}^{-1} (m_D)^T \quad (20)$$

are given by

$$m_{\nu_{1..3}} = 0, (M_1^\nu)^{-1} \lambda^{m+n} (n_1 \mp \sqrt{n_2}) \quad (21)$$

³ With the choice of n, m, q_{θ_ν} above, there would be the possibility of a $O(\lambda M)$ entry in place of the zero for $(M_{\nu^c})_{22}$; this does not change any of our conclusions in an essential way (in particular the matrix m_ν below retains a zero eigenvalue).

where

$$\begin{aligned} n_1 &= a_\nu d_\nu + b_\nu e_\nu + c_\nu f_\nu, \\ n_2 &= d_\nu^2 a_\nu^2 + e_\nu^2 b_\nu^2 + f_\nu^2 c_\nu^2 + e_\nu^2 c_\nu^2 + b_\nu^2 f_\nu^2 + d_\nu^2 b_\nu^2 + a_\nu^2 e_\nu^2 + d_\nu^2 c_\nu^2 + a_\nu^2 f_\nu^2. \end{aligned} \quad (22)$$

It is clear that $|m_{\nu_3}| > |m_{\nu_2}|$ and although both are the same order in λ the relation $|m_{\nu_3}| = 10|m_{\nu_2}|$ holds for “reasonable values” of $a_\nu \cdots f_\nu$. Such a hierarchy accommodates solar and atmospheric neutrino data-see for example Ref. [41]. With masses⁴

$$m_{\nu_1} = 0, \quad m_{\nu_2} = 5 \times 10^{-3} \text{eV}, \quad m_{\nu_3} = 5 \times 10^{-2} \text{eV} \quad (23)$$

we would also expect large mixing angles $\theta_{12}^\nu, \theta_{23}^\nu$ and a small mixing angle θ_{13}^ν in the rotation to the neutrino mass eigenstate basis from the charged lepton mass eigenstate basis. It is easy to construct examples (without fine-tuning) that give rise to precisely this structure within our scenario. In order to generate a neutrino spectrum in the region of Eq. (23), we would require (assuming $a_\nu \cdots f_\nu$ are $O(1)$)

$$M_1^\nu \sim \lambda^{(m+n)} 10^{16} \text{GeV}, \quad (24)$$

or $M_1^\nu \sim 10^{14} \text{GeV}$ in the case $m+n=3$. An example consistent with our requirements is given by:

$$m_L = \begin{pmatrix} -0.56\lambda^4 & 0.56\lambda^2 & -1.07 \\ -0.36\lambda^4 & -2.11\lambda^2 & -0.22 \\ -0.49\lambda^4 & -0.14\lambda^2 & -1.43 \end{pmatrix}, \quad m_D = \begin{pmatrix} \lambda^2 & \lambda \\ 2\lambda^2 & 5\lambda \\ 3\lambda^2 & 1.9\lambda \end{pmatrix} v_2, \quad (25)$$

(where m_L is the charged lepton mass matrix) when we obtain the neutrino spectrum of Eq. (23) for $M_1^\nu \sim 2.4 \times 10^{14} \text{GeV}$, with $\theta_{12}^\nu = 0.53$, $\theta_{23}^\nu = 0.78$ and $\theta_{13}^\nu = 0$. Of course the result for M_1^ν is sensitive to the overall scale of m_D .

We have therefore achieved our goal of incorporating neutrino masses into the AMSB paradigm; with the added bonus that no additional symmetries (beyond U_1') are required to adequately suppress proton decay. The chief feature distinguishing the sparticle spectrum from the massless neutrino case considered in Ref. [13] is the splitting between the first and second generation right-handed squarks and sleptons (the degeneracy persists in the LH case because of the generation independent doublet U_1' charges).

⁴ The recently reported measurement [42] in neutrinoless double beta decay of a Majorana neutrino mass in the region $0.11 - 0.56 \text{eV}$ is not readily accommodated within our scenario; we note, however, some controversy[43] regarding this result.

For example, with $\tan\beta = 5$, gravitino mass $m_0 = 40\text{TeV}$, $\zeta_1 = -0.02\text{TeV}^2$ and $\zeta_2 = 0.0227\text{TeV}^2$, $h_2 = 12$, $t_1^c = -7/2$, we find $|\mu| = 571\text{GeV}$, and choosing $\text{sign}\mu = -1$ we obtain the following spectrum:

$$\begin{aligned}
m_{\tilde{t}_1} &= 869, & m_{\tilde{t}_2} &= 484, & m_{\tilde{b}_1} &= 825, & m_{\tilde{b}_2} &= 1082, & m_{\tilde{\tau}_1} &= 148, \\
m_{\tilde{\tau}_2} &= 442, & m_{\tilde{u}_L, \tilde{c}_L} &= 931, & m_{\tilde{u}_R} &= 908, & m_{\tilde{c}_R} &= 856 & m_{\tilde{d}_L, \tilde{s}_L} &= 934, \\
m_{\tilde{d}_R} &= 998, & m_{\tilde{s}_R} &= 1042, & m_{\tilde{e}_L, \tilde{\mu}_L} &= 149, & m_{\tilde{e}_R} &= 117, \\
m_{\tilde{\mu}_R} &= 323 & m_{\tilde{\nu}_e, \tilde{\nu}_\mu} &= 126, & m_{\tilde{\nu}_\tau} &= 125 & m_{h, H} &= 122, 166, \\
m_A &= 161, & m_{H^\pm} &= 181, & m_{\tilde{\chi}_{1,2}^\pm} &= 112, 575, \\
m_{\tilde{\chi}_{1\dots 4}} &= 111, 369, 579, 579 & m_{\tilde{g}} &= 1007,
\end{aligned} \tag{26}$$

where all masses are given in GeV. The squarks \tilde{t}_1, \tilde{b}_1 and $\tilde{\tau}_1$ couple more strongly to t_L, b_L and τ_L respectively, though (for our chosen $\tan\beta$) the $\tilde{t}_{1,2}$ mixing is of course substantial. For a given m_0 , an acceptable vacuum is obtainable for only a very restricted range of the parameters $\tan\beta, \zeta_1, \zeta_2$. The main constraints on the boundary of the allowed region come from the slepton and Higgs masses. The triangular region in the $\zeta_{1,2}$ plane which meets these constraints is shown in Fig. 1 (for $m_0 = 40\text{TeV}$, and $\tan\beta = 5$). The LSP can be the neutral wino (as in Eq. (26) above), or the $\tilde{\nu}_\tau$; note, however, that there is a small region in which the \tilde{e}_R is the LSP, which we would exclude on cosmological grounds. The point corresponding to our example of Eq. (26) is indicated by an asterisk in the diagram.

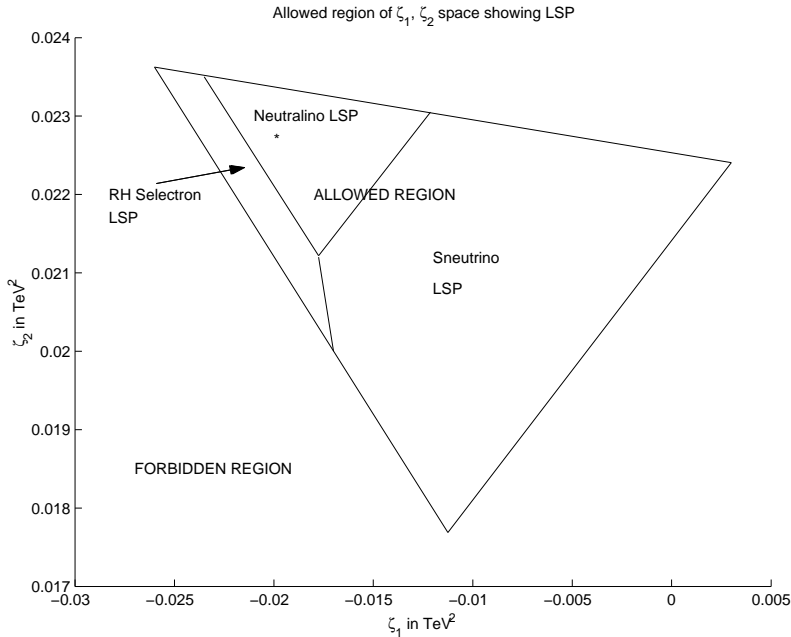


Fig.1: Allowed values of $\zeta_{1,2}$ for $\tan\beta = 5$, $m_0 = 40\text{TeV}$ and $\text{sign}\mu = -1$.

For discussion of sneutrino cold dark matter, see Refs. [44]. For $m_0 \sim 40 \rightarrow 50\text{TeV}$ we find a maximum possible value of $\tan\beta \sim 15$. With, for example, $m_0 = 50\text{TeV}$ and $\tan\beta = 10$ we find for $\zeta_1 = -0.03, \zeta_2 = 0.032$ a spectrum similar to the above with generally increased masses. The spectrum always features the near-degenerate wino triplet that is characteristic of the AMSB scenario. However when the LSP is the $\tilde{\nu}_\tau$, the dominant decay modes of the charged and neutral winos will be to $\tilde{\nu}_\tau$ accompanied in the first case by a charged lepton and in the second by a neutrino. There is substantial $\tilde{t}_{L,R}$ mixing and consequently radiative corrections raise m_h above the current bound $m_h > 114\text{GeV}$ [45] (we have included explicit radiative corrections other than leading logarithm effects in the calculation of m_h only). Notice that $\tilde{e}_L, \tilde{\mu}_L$ are quite light; in the so-called mAMSB model (where the slepton mass problem is resolved by adding a common $(\text{mass})^2$ to the scalars) this would be disfavoured due to the existence of charge-breaking extrema of the scalar potential[29]. We will investigate this possibility in our context elsewhere.

We must consider the issue of scalar-mediated FCNCs, which at first sight would appear to pose a real problem, because of our generation-dependent U'_1 charges. The AMSB contributions to the scalar masses are diagonalised to a good approximation when we transform to the fermion mass-diagonal basis; as in fact are the FI contributions to the LH squarks and sleptons, because of the universal doublet U'_1 charges. Thus the main source of supersymmetric FCNC effects is potentially from the RH squarks and sleptons. In the case of the squarks, these effects can be reduced by increasing the gravitino mass m_0 , which determines the scale of the AMSB contributions. However in the case of the sleptons, because of the crucial role of the FI terms, it is generally the case that some of the sleptons are comparatively light. In the sample spectrum (Eq. (26)) note in particular that $m_{\tilde{e}_R} = 117\text{GeV}$. As we indicated above, what in fact saves us from trouble is the fact that our choice of texture matrices means that the LH charged lepton masses are diagonalised by transforming (to a good approximation) the *LH fields only*.

Consider, for example, the contribution to $\mu \rightarrow e\gamma$ from the neutralino/RH charged slepton loop. Because this contribution to the branching ratio is suppressed compared to the chargino/LH sneutrino loop we are able to tolerate a larger amount of flavour mixing than when this mixing occurs in the LH sector. We find, in fact, that we typically obtain

$$\delta_{e\mu}^{RR} = \frac{m_{\tilde{e}_R\tilde{\mu}_R}^2}{m_{\tilde{e}_R}^2} \sim 10^{-2} \quad (27)$$

and that this leads to sufficient suppression of the branching ratio for $\mu \rightarrow e\gamma$ for the kind of spectrum shown in Eq. (26).

The correlation between the supersymmetric contributions to $(g - 2)_\mu$ and $\mu \rightarrow e\gamma$ that has been discussed by a number of authors[40][46] is weakened here because the former is dominated by the chargino/LH sneutrino loop. For choices of ζ_1, ζ_2 such that $\tilde{\mu}_L, \tilde{e}_L$ are light compared to $\tilde{\mu}_R, \tilde{e}_R$ it is possible to obtain within our framework a supersymmetric contribution to $(g - 2)_\mu$ sufficient to explain the (now reduced [47]) discrepancy between the Standard Model and experiment[48] while being consistent with the limit on $\mu \rightarrow e\gamma$.

We now turn to the $\Delta \neq 0$ case. The Higgs μ term is then forbidden by the U'_1 symmetry, but if $|\Delta|$ is sufficiently large, then we can imagine generating the μ -term via the FN-mechanism[49], i.e. via an interaction of the form $M_U H_1 H_2 (\frac{\theta_H}{M_U})^{n_H}$ where $n_H \approx 16$. The resulting sparticle spectrum is broadly similar to that considered above in the $\Delta = 0$ case, and neutrino masses can be introduced in a similar way. Depending on the θ -charge assignments, it may be that whereas proton decay is adequately suppressed, lepton number-violating R -parity-violating operators are generated by the texture mechanism causing decay of the LSP.

To summarise: with a generation-dependent U'_1 charge assignment we are able to extend our previous FI-based solution to the AMSB tachyonic slepton problem to accommodate neutrino masses. We achieve natural suppression of leptonic FCNCs via a combination of universal doublet U'_1 charges and textures which are diagonalised primarily by LH transformations. We will consider in more detail elsewhere the constraints placed on our general framework by experimental limits on sparticle masses, and other issues such as vacuum stability and CP-violation.

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