Published for SISSA by Springer

Received: April 17, 2012
Accepted: May 26, 2012
Published: June 21, 2012

Flavored gauge-mediation

Yael Shadmi and Peter Z. Szabo

Physics Department, Technion-Israel Institute of Technology,
Haifa 32000, Israel

E-mail: yshadmi@physics.technion.ac.il, peter.z.szabo@gmail.com

Abstract: The messengers of Gauge-Mediation Models can couple to standard-model matter fields through renormalizable superpotential couplings. These matter-messenger couplings generate generation-dependent sfermion masses and are therefore usually forbidden by discrete symmetries. However, the non-trivial structure of the standard-model Yukawa couplings hints at some underlying flavor theory, which would necessarily control the sizes of the matter-messenger couplings as well. Thus for example, if the doublet messenger and the Higgs have the same properties under the flavor theory, the resulting messenger-lepton couplings are parametrically of the same order as the lepton Yukawas, so that slepton mass-splittings are similar to those of minimally-flavor-violating models and therefore satisfy bounds on flavor-violation, with, however, slepton mixings that are potentially large. Assuming that fermion masses are explained by a flavor symmetry, we construct viable and natural models with messenger-lepton couplings controlled by the flavor symmetry. The resulting slepton spectra are unusual and interesting, with slepton mass-splittings and mixings that may be probed at the LHC. In particular, since the new contributions are typically negative, and since they are often larger for the first- and second-generation sleptons, some of these examples have the selectron or the smuon as the lightest slepton, with mass splittings of a few to tens of GeV.

Keywords: Supersymmetry Phenomenology

ArXiv ePrint: 1103.0292

doi:10.1007/JHEP06(2012)124

Contents

1 Introduction

Motivated by the absence of flavor changing neutral currents and rare decays, most studies of supersymmetry at colliders assume universal sfermion masses at the scale where supersymmetry breaking is mediated to the Minimal Supersymmetric Standard Model (MSSM). Any sfermion mass splittings or mixings then originate from the Standard Model (SM) Yukawa couplings only, and are negligibly small, except for the stop mass splitting and, for large $\tan\beta$, also the sbottom and stau mass splittings. Such models, in which the SM Yukawas are the only source of generation-dependence, are usually referred to as Minimally Flavor Violating (MFV). The assumption of MFV is too restrictive however. Current constraints on lepton-violating decays [1, 2] for example allow for slepton mass splittings and mixings that may well be observable at the LHC (see for example [3]). If such splittings and mixings are indeed observed, they would provide a wealth of information about the origin of supersymmetry breaking, and quite possibly, about the origin of the SM fermion masses.

It is interesting to ask therefore if there are viable models of supersymmetry breaking that give rise to appreciable departures from sfermion mass universality, and several classes of models were recently discussed in the literature [3–7]. Here we will present another example which is particularly simple, namely, Minimal Gauge Mediated Supersymmetry Breaking (GMSB) Models [8, 9] with messenger-matter couplings.

The main appeal of GMSB models of course is that the soft masses are generated by gauge interactions and are therefore generation-independent by construction. In practice, however, in the most successful examples of GMSB, the soft masses are generated by loops of messenger fields with SM gauge quantum numbers, which can have renormalizable superpotential couplings to the MSSM [10–13]. Such couplings would lead to generation-dependent sfermion masses and flavor changing neutral currents, so one usually invokes some global symmetries in order to forbid them. Here we will take a more liberal approach towards messenger-matter couplings, and show that they can result in viable models with rich and interesting spectra.[1]

Consider for example the standard set of vectorlike $5+\bar{5}$ messengers [8, 9]. Of these, one of the SU(2) doublets, which we will denote by D, is in the same SM representation as the down-type Higgs, H_D. Assuming that it has the same R-parity assignment as the Higgs, the superpotential can contain terms of the form

$$y_L D\, l\, e\,, \tag{1.1}$$

in addition to the usual Yukawa

$$Y_L H_D\, l\, e\ . \tag{1.2}$$

Here l is the lepton doublet, e is the lepton singlet, and Y_L and y_L are 3×3 matrices of couplings. For an arbitrary matrix y_L, one would have disastrous flavor changing processes. However, the SM Yukawa matrix Y_L is far from arbitrary. Most of its entries are very small, hinting at some underlying flavor theory. So it is not implausible that the same underlying theory would also suppress the entries of the new coupling y_L, so that the two matrices are of the same order of magnitude. All flavor constraints would then be satisfied, because the model is qualitatively MFV: all generation dependence originates from couplings which, albeit new, are of the same order of magnitude as the SM Yukawas, and the resulting mass splittings are similar to those of MFV models. Even this minimal scenario has interesting phenomenological implications. Since the two matrices Y_L and y_L are not proportional to each other, slepton mixings can be appreciable. As a concrete example, assume that the Yukawa matrix is governed by an abelian flavor (Froggatt-Nielsen) [17] symmetry. If D and H_D carry the same charge under the flavor symmetry, each entry of the matrix y_L is parametrically the same as the corresponding entry in the matrix Y_L, realizing the minimal scenario described above. As we will see, one can also construct models in which y_L is very different from Y_L, leading to large splittings between the first two generations, and with the selectron or smuon being lighter than the stau.

All our models are GMSB models with matter-messenger couplings controlled by the same flavor symmetry which generates the structure of fermion mass matrices. Since the slepton masses, and in particular, the selectron and smuon masses, will probably be the easiest probes of flavor dependence at the LHC, we focus on models in which the only new messenger couplings involve leptons. The LHC signatures of generation dependent slepton spectra have received a lot of attention recently (see for example [18–33]). It would be interesting to generalize our results to the squarks as well.

In fact, the largest couplings are often the couplings to the first generations, so the selectron or smuon exhibit the largest mass splitting. The reason is that, in order to obtain

[1] Matter-messenger couplings were also studied in the context of triplet seesaw models, with the messengers in a $15+\bar{15}$ of SU(5) [14–16].

JHEP06(2012)124

appreciable mass splittings, we need some entries of y_L to be larger than the corresponding entry in Y_L. This happens if the flavor charge of D is smaller than the flavor charge of H_D (adopting the standard convention that the flavor spurion has negative charge). On the other hand, the third-generation fields must have smaller flavor charges than the first- and second-generation fields, in order for their masses to be less suppressed. Generically then, the entries of y_L corresponding to the third generation have overall negative charge, and since the superpotential can only contain positive powers of the spurion, cannot appear in the superpotential.

The couplings of eq. (1.1) were studied in [13], motivated by the fact that they mediate messenger decay, and thus solve the cosmological problems associated with stable messengers. Unlike our models, the models of [13] were MFV, with the messengers and SM living on different branes in a 5d setup so that the couplings eq. (1.1) originate solely from Higgs-messenger mixings, with $y = \epsilon Y$ and with ϵ suppressed by the size of the extra-dimension.[2]

We will classify the different possible messenger-matter couplings in section 2, and present the basic superpotential of our models in section 3. In section 4 we discuss a few example models and their spectra.

2 General matter-messenger couplings

Minimal GMSB models [8, 9] involve N_5 pairs of vector-like messengers transforming as $5 + \bar{5}$'s of SU(5), coupled to a SM gauge singlet X, whose vacuum expectation value (VEV) $\langle X \rangle \equiv M$ gives mass to the messengers, and whose F-term is non-zero, leading to supersymmetry-breaking splittings in the messenger spectrum. Under the SM gauge group, the messengers transform as

$$T_I \sim (3,1)_{-1/3} \quad \bar{T}_I \sim (\bar{3},1)_{1/3} \quad D_I \sim (1,2)_{-1/2} \quad \bar{D}_I \sim (1,2)_{1/2}\,, \tag{2.1}$$

where $I = 1 \ldots N_5$. For ease of notation, we will define in the following

$$D \equiv D_1\ . \tag{2.2}$$

The possible trilinear superpotential couplings of the messengers to the SM depend on the messengers R-parity charge assignment. For R-parity odd messengers, the most general trilinear superpotential is of the form [11, 12]

$$W_{\rm odd} = H_D q T + H_D D e^c, \tag{2.3}$$

where q denotes the doublet quarks, and e^c denotes the singlet leptons. This superpotential breaks baryon- and lepton-number. For R-parity even messengers one can have the doublet-messenger couplings [12, 13],

$$W^D_{\rm even} = y_U \bar{D} q u^c + y_D D q d^c + y_L D l e^c, \tag{2.4}$$

[2]This can be achieved in 4d too, using some broken global symmetry to distinguish between D and H_D, with the suppression factor ϵ being the relevant spurion.

where u^c and d^c are the singlet up and down quarks respectively, l is the lepton doublet, and the y's are 3×3 matrices of couplings, (throughout, we use small letters for matter fields to distinguish them from the Higgses and messengers, which we denote by capital letters), as well as the triplet-messenger couplings

$$W^T_{\text{even}} = Tqq + Tu^c e^c + \bar{T}ql + \bar{T}u^c d^c. \tag{2.5}$$

The couplings of eq. (2.5) are precisely the same as the couplings of GUT triplet Higgses, which have the same SM charges as T and $\bar{T}$, and would lead to proton decay. Here we will assume that the messengers have the same R-parity charge assignment as the Higgses, and impose an additional Z_2 symmetry which forbids the triplet couplings eq. (2.5). The relevant superpotential is then eq. (2.4).

The couplings of eqs. (2.3)–(2.5) generate sfermion masses squared starting at one-loop [11, 12], but the one-loop contributions vanish at leading order in the supersymmetry breaking, so that in the limit of small supersymmetry breaking, the dominant contributions are the two-loop analogs of the usual gauge contributions. Here we will concentrate on these two-loop contributions. Unlike in minimal GMSB models, the new couplings also generate A-terms at one-loop. The dependence of the soft terms on the matrices Y and y can be inferred from a spurion analysis as in [34], treating Y and y as spurions of the SM $SU(3)^5$ flavor symmetry. Since the abelian Froggatt-Nielsen symmetry that we will invoke in the following only determines the matrices Y and y up to $O(1)$ coefficients, such a spurion analysis is completely adequate for our purposes. Still, we will explicitly compute the mixed gauge-Yukawa contributions to the soft terms.[3] As we will see, the gauge-Yukawa contributions will be the dominant contributions in our models, and knowing their signs will allow us to determine the hierarchy in the slepton spectrum.

Since we are mainly interested in the implications for the slepton spectrum, our models are constructed so that y_U always vanishes and y_D is negligible (or zero). The slepton masses are then,

$$m^2_{\tilde{l}} = \frac{1}{128\pi^2}\left[N_5\left(\frac{3}{4}g_2^4 + \frac{5}{3}g_Y^4\right)\mathbf{1} - \left(\frac{3}{2}g_2^2 + 6g_Y^2\right)y_L y_L^\dagger + \ldots\right]\left|\frac{F}{M}\right|^2, \tag{2.6}$$

$$m^2_{\tilde{e}^c} = \frac{1}{128\pi^2}\left[N_5\left(\frac{20}{3}g_Y^4\right)\mathbf{1} - \left(3g_2^2 + 12g_Y^2\right)y_L^\dagger y_L + \ldots\right]\left|\frac{F}{M}\right|^2. \tag{2.7}$$

The first terms in eq. (2.6), eq. (2.7) are the usual GMSB contributions, which are proportional to the number of messenger pairs N_5. The remaining terms are new contributions and lead to mass splittings and mixings among the different generations. The latter can be appreciable even for small y_L's, since the GMSB contribution to the soft mass is proportional to the identity matrix [3]. The ellipses stand for pure Yukawa terms including terms with four powers of the matrices y, and terms with two powers of y and two powers of Y, such as $y_L y_L^\dagger y_L y_L^\dagger$, $y_L y_L^\dagger Y_L Y_L^\dagger$ + h.c.. Up to order one coefficients, these terms can

[3] We derive these using the method of [12], generalizing the results of [13] to the case of 3-generations, since we are particularly interested in large couplings of the messengers to the first and second generation scalars.

Superfield	R-parity	Z_3	Z_2
X	even	1	even
T_1	even	0	odd
$\bar{T}_1$	even	-1	odd
D	even	0	even
$\bar{D}_1$	even	-1	even
$T_I, \bar{T}_I, D_I, \bar{D}_I$ $(I = 2, \ldots, N_5)$	even	1	even
q, u^c, d^c, l, e^c	odd	0	even
H_U, H_D	even	0	even

Table 1. $Z_3 \times Z_2$ symmetry charges.

be determined by an SU(3)^5 spurion analysis [34]. In all of our models, the pure Yukawa terms are negligible compared to the mixed gauge-Yukawa terms, so we can safely ignore them.[4] The A terms are given by,

$$A_L = -\frac{1}{16\pi^2}\left[y_L y_L^\dagger Y_L + 2Y_L y_L^\dagger y_L\right]\frac{F}{M}\,. \tag{2.8}$$

We note that the structure of our models is similar to the gauge-gravity hybrid models of [3], in which the universal contribution is also gauge-mediated, with a gravity-mediated generation-dependent contribution which is important for a high messenger scale. In both frameworks, the size of the non-universal contribution is controlled by a flavor symmetry, and flavor constraints are satisfied through the interplay of degeneracy and alignment [35].

3 Basic superpotential

In addition to R-parity, we will impose a $Z_3 \times Z_2$ symmetry on the theory, with charges given in table 1.[5] The most general superpotential allowed by this symmetry is

$$\begin{aligned} W =&\; X\left(XX + T_I\bar{T}_I + D_I\bar{D}_I + H_D\bar{D}_1\right) + H_U q u^c + H_D q d^c + H_D l e^c \\ &+ D q d^c + D l e^c\,, \end{aligned} \tag{3.1}$$

where we omitted the generalized μ-terms, $H_U H_D + H_U D$. Just like the usual μ-term, these can be forbidden by some Peccei-Quinn symmetry, and we will not consider them in the following.

The first line of eq. (3.1) contains the messenger couplings to the supersymmetry-breaking sector as well as the usual Yukawa terms. We explicitly display here the term X^3, which is typically needed in order to generate appropriate VEVs for X, and motivates our choice of a Z_3 symmetry. We will not consider this term further.

The second line of eq. (3.1) is our focus here, with the messenger field D replacing H_D. The analogous up-type messenger-matter coupling $\bar{D}qu^c$ is eliminated by the Z_3 symmetry.

[4]We thank Anna Rossi for pointing out to us an error in some of the pure Yukawa terms in an earlier version of this paper.

[5]We note that these models, while preserving coupling unification, are incompatible with grand unification, since the triplet- and doublet-messengers carry different $Z_3 \times Z_2$ charges.

JHEP06(2012)124

It is simple to allow for this term as well. To do so, one must use at least two separate pairs of messengers, the first charged as shown in table 1 for $I = 1$, and the second, with the charges of D and $\bar{D}$ swapped. Since we are interested in slepton masses here, it is simplest to stick to the charges of table 1, so that the new couplings only involve the leptons and down-quarks. As we will see later, it is often possible to impose additional symmetries on the models so that down-quark couplings are eliminated as well.

Note that, in this construction, the new couplings of the messengers to down quarks and leptons (or alternatively, to up-quarks) can appear with one set of messengers, $N_5 = 1$. Having both up-type and down-type messenger couplings requires however $N_5 > 1$.

3.1 MFV-like masses

So far, H_D and D have the same charges under all the symmetries of the model. If this remains true in the presence of any additional symmetries, we can define D as the combination of D and H_D that couples to X, and take H_D to be the orthogonal combination. The superpotential eq. (3.1) then takes the form

$$\begin{aligned} W = X\left(XX + T_I\bar{T}_I + D_I\bar{D}_I\right) + Y_U H_U q u^c \;\; &+ Y_D H_D q d^c + Y_L H_D l e^c \\ &+ y_D D q d^c + y_L D l e^c\,, \end{aligned} \tag{3.2}$$

where we display also the 3×3 matrices of couplings, with Y_U, Y_D and Y_L denoting the usual up-, down-, and lepton-Yukawas respectively, and y_D and y_L denoting the corresponding new couplings. In this case, if the Yuakwa matrices are controlled by some underlying theory, then the matrices y_L and Y_L (and similarly, y_D and Y_D) are parametrically the same. These models are therefore quite similar to MFV models. They contain new matrices of couplings, which, while not proportional to the Yukawa matrices, satisfy

$$(y_L)_{ij} = c_{ij}\,(Y_L)_{ij}\,, \tag{3.3}$$

where c_{ij} are order-one coefficients and $i, j = 1, 2, 3$ are generation indices. The resulting mass-splittings are therefore of the same order of magnitude as those obtained in MFV models. In particular, the first and second generation scalars are practically degenerate. As we will see in section 4.1, such slepton mass splittings are consistent with bounds on rare-decays even for large mixings. On the other hand, the inter-generational mixings are model dependent, and can be large. The masses of down squarks are more stringently constrained by bounds on flavor-changing processes, but still, at least for small $\tan\beta$, the resulting y_D couplings are viable. Here too, one can construct models with large down-quark mixings, but we leave the phenomenology of such models for future study.

The model of section 4.1 provides a concrete realization of eq. (3.3) using a flavor symmetry, but the approximate equality of the messenger couplings and the Yukwas can hold much more generally whenever the messengers and the Higgses have the same properties with respect to the underlying theory of flavor.

3.2 New mass patterns

It is also possible to construct models with additional symmetries, under which D and H_D transform differently. Most of the models we consider below are of this type. In all of these, the term $XH_D D$ is forbidden by holomorphy, so that the superpotential is again of the form eq. (3.2). To illustrate the basic mechanism consider a one-generation toy model. We impose a U(1) symmetry broken by a spurion ϵ of charge -1, with the following charges,

$$H_D\ (-1), \quad d^c\ (1), \quad e^c\ (1), \quad l\ (n \geq 0)\,, \tag{3.4}$$

and all other fields neutral. The term $XH_D\bar{D}_1$ cannot appear while the usual Yukawas are allowed. In this case, the coupling y_L is smaller than Y_L by the factor ϵ, and the phenomenology of this model is not very interesting because the deviations from GMSB masses would be smaller than those induced by the Yukawas. In the models we construct below, however, the new U(1) will be part of a U(1)$\times$U(1) flavor symmetry, with the second U(1) factor compensating for this suppression, and leading to some entries $(y_L)_{ij} > (Y_L)_{ij}$.

4 Generation dependent slepton spectra with a flavor symmetry

We will assume that the hierarchies of the SM fermion masses are explained by a broken flavor symmetry, which we take to be $\mathrm{U}(1)_1\times\mathrm{U}(1)_2$, with each U(1) factor broken by a spurion $\lambda_{1,2}$ of charge -1, and with $\lambda_1 \sim \lambda_2 = \lambda \sim 0.1-0.2$.

The models are then completely specified by choosing $\mathrm{U}(1)_1 \times \mathrm{U}(1)_2$ charges for the different fields. We always take H_U, as well as all the messengers apart from $D \equiv D_1$ to be neutral under this symmetry. In addition, we choose the charges of H_D as $(0,-1)$, with the -1 motivated by the fact that we want to eliminate H_D couplings to the supersymmetry-breaking field X as explained in section 3.2. In fact, the $\mathrm{U}(1)_2$ factor plays the role of the U(1) symmetry of the toy model of that section. The models thus differ from each other because of the charges of the matter fields and the messenger D, and we will discuss different options below.

Since the SM matter fields transform non-trivially under the flavor symmetry, the structure of the new coupling matrices y_L is affected by this symmetry as well, with some entries suppressed by powers of λ, so that the flavor-changing contributions are potentially suppressed by powers of λ.

In order to estimate these contributions and to determine whether the models are viable, it is useful to work in terms of the quantities $\delta_{i\neq j}$ [36], which are the basic quantities constrained by bounds on flavor violation. Since we will be interested in the phenomenological predictions of the models, it is useful to work in the slepton-mass basis, so that the slepton mass differences and mixings are transparent. One can then write (see for example, [37]),[6]

$$\delta^A_{ij} \equiv \frac{\Delta\tilde{M}^2_{Aji}}{\tilde{M}^2_{Aji}}\, K^A_{ij}\,, \tag{4.1}$$

where $A = L$ ($A = R$) refers to the lepton doublets (singlets),

$$\begin{aligned} \Delta\tilde{M}^2_{Aji} &= \tilde{M}^2_{Aj} - \tilde{M}^2_{Ai}\,, \\ \tilde{M}_{Aji} &= \left[\tilde{M}_{Aj} + \tilde{M}_{Ai}\right]/2\,, \end{aligned} \tag{4.2}$$

[6] Neglecting LR mixings, which is a good approximation in the models below.

Superfield	l_1	l_2	l_3	e^c_1	e^c_2	e^c_3	H_d	D
$\mathrm{U}(1)_1$	4	2	0	1	1	0	0	0
$\mathrm{U}(1)_2$	0	2	4	1	−1	−2	−1	−1

Table 2. Flavor charges of MFV-like model.

and where M_{Ai} is the mass of the slepton i, and K^A is the mixing matrix of the electroweak gaugino couplings.[7] Clearly, the flavor-changing contributions can be small if either the mass-splittings or the inter-generation mixings are small, or both. The example below will interpolate between these options.

It will be convenient for our purposes to parametrize the experimental bounds as powers of λ. The most stringent bounds are from [1, 2], and using the results of [38], we have

$$\begin{array}{lll} \delta^L_{12} \lesssim \lambda^4 \,, & \delta^L_{13} \lesssim \lambda - \lambda^2 \,, & \delta^L_{23} \lesssim \lambda \,, \\ \delta^R_{12} \lesssim \lambda^2 \,, & \delta^R_{13} \lesssim \lambda \,, & \delta^R_{23} \lesssim \lambda \,, \\ \delta^{LR}_{12} \lesssim \lambda^5 \,, & \delta^{LR}_{13} \lesssim \lambda^2 \,, & \delta^{LR}_{23} \lesssim \lambda^2 \,. \end{array} \tag{4.3}$$

4.1 MFV-like masses with potentially large mixings

Choosing D and H_D to have identical flavor charges results in MFV-like masses, since y_L and Y_L are equal up to $O(1)$ coefficients. With the flavor charges given in table 2, the desired lepton masses are obtained, and

$$y_L \sim Y_L \sim \begin{pmatrix} \lambda^5 & 0 & 0 \\ \lambda^5 & \lambda^3 & 0 \\ \lambda^5 & \lambda^3 & \lambda \end{pmatrix} \,. \tag{4.4}$$

Here and in the following, the entries are determined to leading order in λ and up to $\mathcal{O}(1)$ coefficients.

One then finds, setting all terms suppressed by more than six powers of λ to zero (we denote such terms by "~ 0"),

$$\widetilde{m}^2_{LL} \sim \frac{\Lambda^2}{128\pi^4} \left[N_5\, G_L \mathbf{1_{3\times 3}} - \frac{3}{2} G_1 \begin{pmatrix} \sim 0 & \sim 0 & \sim 0 \\ \sim 0 & \lambda^6 & \lambda^6 \\ \sim 0 & \lambda^6 & \lambda^2 \end{pmatrix} \right] \tag{4.5}$$

and

$$\widetilde{m}^2_{RR} \sim \frac{\Lambda^2}{128\pi^4} \left[N_5\, G_R \mathbf{1_{3\times 3}} - 3\, G_1 \begin{pmatrix} \sim 0 & \sim 0 & \lambda^6 \\ \sim 0 & \lambda^6 & \lambda^4 \\ \lambda^6 & \lambda^4 & \lambda^2 \end{pmatrix} \right] , \tag{4.6}$$

where we defined the mass scale $\Lambda \equiv F/M$ and the dimensionless numbers

$$G_L \equiv \frac{3}{4} g_2^4 + \frac{5}{3} g_Y^4 \,, \qquad G_R \equiv \frac{20}{3} g_Y^4 \,, \qquad G_1 \equiv g_2^2 + 4 g_Y^2 \,. \tag{4.7}$$

[7] With a slight abuse of notation, we use the same indices to label lepton and slepton states.

The first term of each mass matrix in eq. (4.5), eq. (4.6) is the ordinary GMSB result and the second term is the contribution due to the new messenger-matter couplings. Note that the signs of the diagonal entries in these new contributions are known: The $O(1)$ numbers multiplying the powers of λ on the diagonals are positive. We also get

$$\widetilde{m}_{LR}^2 \sim -\frac{\Lambda v_d}{16\pi^2}\left[3\begin{pmatrix}\sim 0 & \sim 0 & \sim 0\\ \sim 0 & \sim 0 & \sim 0\\ \sim 0 & \lambda^5 & \lambda^3\end{pmatrix} + \frac{\mu}{\Lambda/16\pi^2}\tan\beta\begin{pmatrix}\lambda^5 & 0 & 0\\ \lambda^5 & \lambda^3 & 0\\ \lambda^5 & \lambda^3 & \lambda\end{pmatrix}\right], \tag{4.8}$$

The second term of $\widetilde{m}_{LR}^2$ is the standard μ-term contribution, while the first comes from the A-term, and is sub-dominant even for $\tan\beta \sim 1$.

Since the mass splittings in this case are of the order of the mass splittings in MFV models, the model automatically satisfies all flavor constraints, with a selectron and smuon that are practically degenerate. The stau mass is split from the other masses by $O(\lambda^2)$, coming from the 3-3 entries of the LL, RR and LR blocks (the latter appears in minimal GMSB models too). A similar effect is induced by the running from the messenger scale to the weak scale, and is probably the dominant effect since it's log-enhanced. This too is a feature of minimal GMSB models, so the stau splitting here is the same as in GMSB models, and this holds in all of our models.

However, unlike MFV models in which the fermion mass matrix and the slepton mass matrix are diagonal in the same basis, this model predicts $O(\lambda^2)$ mixings of $\tilde{e}_R - \tilde{\mu}_R$ and $\tilde{\mu}_R - \tilde{\tau}_R$. The former might not be observable because of the small selectron-smuon mass splitting, but the latter may be within reach of LHC experiments.

As explained before, this model will necessarily contain couplings of the D messenger to down quarks, so that down squarks receive generation-dependent corrections as well. These are largest when the third generation is involved, with

$$\frac{\Delta\tilde{M}_{ij}^2}{\tilde{M}_{ji}^2} \lesssim \frac{1}{N_5}\, y_b^2\,, \tag{4.9}$$

where y_b is the bottom Yukawa. The most severe constraint on the models is [39] $\delta_{13}^{d,LL}\delta_{13}^{d,RR} \lesssim 5\cdot 10^{-5}$, but this is satisfied for $\tan\beta \sim 1$ or for $N_5 = 3$ even for $\tan\beta \sim 5$.

4.2 Selectron splitting

In order to obtain some large entries in y_L, these entries must involve smaller powers of λ compared to the relevant entry of Y_L. It is easy to achieve this by taking the $\mathrm{U}(1)_1$ charge of D to be smaller than the $\mathrm{U}(1)_1$ charge of H_D (which we took to be zero). Consider for example the flavor charges of table 3. The large negative charge of D has two consequences for the slepton spectrum. First, most of the entries of y_L vanish due to holomorphy [35], with only the $1-1$ entry surviving. Second, this entry is rather large. Thus, only the first-generation fields, whose charges are largest so that their masses would be the most suppressed, couple to the messenger sector, and the modification of the selectron mass is appreciable. In addition, because of this large negative charge, it is easy to choose charges for the down quarks so that y_D vanishes identically.

Superfield	l_1	l_2	l_3	e^c_1	e^c_2	e^c_3	H_d	D
U(1)$_1$	4	2	0	1	1	0	0	−5
U(1)$_2$	0	2	4	1	−1	−2	−1	0

Table 3. Flavor charges for section 4.2.

The resulting lepton Yukawas are as in eq. (4.4) while the new couplings are given by,

$$y_L \sim \begin{pmatrix} \lambda & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} . \tag{4.10}$$

The LL and RR blocks are then,

$$\widetilde{m}^2_{LL} \sim \frac{\Lambda^2}{128\pi^4}\left[N_5\, G_L \mathbf{1_{3\times 3}} - \frac{3}{2} G_1 \begin{pmatrix} \lambda^2 & \sim 0 & \sim 0 \\ \sim 0 & \sim 0 & \sim 0 \\ \sim 0 & \sim 0 & \sim 0 \end{pmatrix}\right] , \tag{4.11}$$

and

$$\widetilde{m}^2_{RR} \sim \frac{\Lambda^2}{128\pi^4}\left[N_5\, G_R \mathbf{1_{3\times 3}} - 3G_1 \begin{pmatrix} \lambda^2 & \sim 0 & \sim 0 \\ \sim 0 & 0 & 0 \\ \sim 0 & 0 & 0 \end{pmatrix}\right] . \tag{4.12}$$

The A-terms are negligible in this model.

The slepton mixings in this case arise solely from the lepton mass matrix, and are given by,

$$K^L_{12} \sim \lambda^4\,, \quad K^L_{13} \sim \lambda^8\,, \quad K^L_{23} \sim \lambda^4\,; \quad K^R_{12} \sim \lambda^2\,, \quad K^R_{13} \sim \lambda^4\,, \quad K^R_{23} \sim \lambda^2\,. \tag{4.13}$$

The only significant δ is $\delta_{RR,12} \sim \lambda^4/N_5$ which is below the bound. In both the L- and the R-sectors, the selectron is lighter than the smuon by $\delta m \sim \lambda^2$. Given that our estimates are parametric only, it is impossible to tell in these models whether the selectron is the lightest slepton, since the stau masses are also driven lower by $O(\lambda^2)$ both by running effects and by the μ term contribution (the RGE contribution could be bigger for a high messenger scale because it is logarithmically enhanced). In any case, the resulting spectrum is very interesting, with the smuon being the heaviest slepton and the selectron and stau lighter than the smuon, with mass splittings around a few GeV or even 10 GeV, and with $e - \mu$ and $\mu - \tau$ mixings of a few percent in the R sector.

4.3 Large mixings

The previous model leads to small mixings of the selectron with the other sleptons. We can also obtain large selectron mixings by choosing charges so that the new couplings are similar for the three generations. In this case, of course, the mass splittings are more constrained, so we want the size of the new couplings to be sufficiently small, motivating the choice of charges for D as shown in table 4.

JHEP06(2012)124

Superfield	l_1	l_2	l_3	e_1^c	e_2^c	e_3^c	H_d	D
$\mathrm{U}(1)_1$	2	2	2	4	2	0	0	0
$\mathrm{U}(1)_2$	0	0	0	1	1	1	-1	0

Table 4. Flavor charges for section 4.3.

This yields an ordinary lepton Yukawa matrix of

$$Y_L \sim \begin{pmatrix} \lambda^6 & \lambda^4 & \lambda^2 \\ \lambda^6 & \lambda^4 & \lambda^2 \\ \lambda^6 & \lambda^4 & \lambda^2 \end{pmatrix} \tag{4.14}$$

which requires a somewhat small $\tan\beta$, and

$$y_L \sim \begin{pmatrix} \lambda^7 & \lambda^5 & \lambda^3 \\ \lambda^7 & \lambda^5 & \lambda^3 \\ \lambda^7 & \lambda^5 & \lambda^3 \end{pmatrix} . \tag{4.15}$$

The resulting slepton mass matrices are then

$$\widetilde{m}^2_{LL} \sim \frac{\Lambda^2}{128\pi^4} \left[N_5\, G_L \mathbf{1_{3\times3}} - \frac{3}{2} G_1 \begin{pmatrix} \lambda^6 & \lambda^6 & \lambda^6 \\ \lambda^6 & \lambda^6 & \lambda^6 \\ \lambda^6 & \lambda^6 & \lambda^6 \end{pmatrix} \right] , \tag{4.16}$$

and

$$\widetilde{m}^2_{RR} \sim \frac{\Lambda^2}{128\pi^4} \left[N_5\, G_R \mathbf{1_{3\times3}} - 3 G_1 \begin{pmatrix} \sim 0 & \sim 0 & \sim 0 \\ \sim 0 & \sim 0 & \sim 0 \\ \sim 0 & \sim 0 & \lambda^6 \end{pmatrix} \right] . \tag{4.17}$$

The A-terms are negligible so that the only contribution to the LR term is the usual μ term contribution. The slepton masses are approximately degenerate in this case, apart from the stau. The mixings of the R-sector are as in eq. (4.13), but the L-sector has $O(1)$ mixings,

$$K^L_{12},\ K^L_{13},\ K^L_{23} \sim O(1) . \tag{4.18}$$

Finally, let us comment on the down sector in this model. With the choice of D charges as in table 4, the down-messenger couplings would generically satisfy

$$y_{D,ij} \sim \lambda Y_{D,ij} . \tag{4.19}$$

Thus, the relative mass splittings in this case are generically $O(\lambda^2)$ smaller than those of eq. (4.9), and the models are consistent with flavor bounds involving down squarks.

4.4 Some large splittings and some large mixings

Finally, we present an example in which the $\tilde{e}_L$ and the $\tilde{\mu}_R$ masses receive significant corrections, with a large $2-3$ mixing in the L-sector. The flavor charges are given in

JHEP06(2012)124

Superfield	l_1	l_2	l_3	e^c_1	e^c_2	e^c_3	H_d	D
$\mathrm{U}(1)_1$	2	2	2	2	2	0	0	-4
$\mathrm{U}(1)_2$	0	2	2	2	0	0	-1	0

Table 5. Flavor charges for section 4.4.

table 5. The lepton Yukawa matrix is as in eq. (4.4), and the messenger-lepton Yukawa couplings are

$$y_L \sim \begin{pmatrix} \lambda^2 & 1 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} . \tag{4.20}$$

Just as in the model of section 4.2, the large and negative charge of D results in a large effect on the second generation, with no effect on the third generation. Furthermore, the messenger couplings to down quarks will also vanish generically, since the total powers of λ that should enter the down mass matrix entries, and therefore the total effective charge of these fields, are typically smaller than those associated with the leptons.

These new couplings lead to

$$\tilde{m}^2_{LL} \sim \frac{\Lambda^2}{128\pi^4}\left[G_L \mathbf{1_{3\times 3}} - \frac{3}{2}G_1 \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix} + \begin{pmatrix} 1 & \sim 0 & \sim 0 \\ \sim 0 & \lambda^6 & \lambda^6 \\ \sim 0 & \lambda^6 & \lambda^6 \end{pmatrix}\right] , \tag{4.21}$$

and

$$\tilde{m}^2_{RR} \sim \frac{\Lambda^2}{128\pi^4}\left[N_5\, G_R \mathbf{1_{3\times 3}} - 3G_1 \begin{pmatrix} \lambda^4 & \lambda^2 & 0 \\ \lambda^2 & 1 & 0 \\ 0 & 0 & 0 \end{pmatrix} + \begin{pmatrix} \lambda^4 & \lambda^2 & \lambda^6 \\ \lambda^2 & 1 & \lambda^4 \\ \lambda^6 & \lambda^4 & 0 \end{pmatrix}\right] . \tag{4.22}$$

The A-terms are again very small, with

$$\tilde{m}^2_{LR} \sim -\frac{\Lambda v_d}{16\pi^2}\left[\begin{pmatrix} \lambda^5 & \sim 0 & 0 \\ \lambda^5 & \lambda^3 & 0 \\ \lambda^5 & \lambda^3 & 0 \end{pmatrix} + \frac{\mu}{\Lambda/16\pi^2}\tan\beta \begin{pmatrix} \lambda^5 & 0 & 0 \\ \lambda^5 & \lambda^3 & \lambda \\ \lambda^5 & \lambda^3 & \lambda \end{pmatrix}\right] . \tag{4.23}$$

The RR mixings are as in eq. (4.13), and the LL mixing are negligible apart from $K^L_{23} = O(1)$. The constrained quantities eq. (4.3) for the LL block are negligible. In the RR block, $\delta_{RR,12}$ and $\delta_{RR,23}$ are of order λ^2/N_5, saturating the bound on $\delta_{RR,12}$ for small N_5. The same holds for $\delta_{LR,12} \sim \lambda^5/N_5$. The other δ_{LR}'s are negligible. Since the model is only specified up to $O(1)$ parameters, we see that it can be consistent with bounds on flavor-violation for parts of the parameter space.

This model has a very interesting spectrum. The $\tilde{e}_L$ has a large mass splitting compared to the other L-sleptons,

$$\frac{\Delta \tilde{M}^2_{L1i}}{\tilde{M}^2_L} \sim \frac{1}{N_5}, \quad i = 2, 3 , \tag{4.24}$$

and hardly mixes with the $\tilde{\mu}_L$, $\tilde{\tau}_L$. In addition, the $\tilde{\mu}_L - \tilde{\tau}_L$ mixing is large.

In the R-sector,

$$\frac{\Delta\tilde{M}^2_{R2i}}{\tilde{M}^2_L} \sim \frac{1}{N_5}\,, \quad i = 1, 3 \tag{4.25}$$

so that the R-smuon is significantly split from the other R-sleptons. Since, in addition, the masses of $\widetilde{e}_R$ and the staus have $O(\lambda^2)$ corrections to their GMSB masses, all six sleptons are separated in mass.

To conclude this section, we comment on the flavor-violating processes mediated by tree-level exchange of the doublet messenger D. Since this messenger field has Yukawa couplings to the matter fields, integrating it out gives rise to higher-dimension terms coming from tree-level exchange of either the scalar- or fermion-D. In principle, such terms can mediate flavor-violationg processes such as $\mu \to eee$ or $\mu \to e\gamma$. However, at leading order, such processes are proportional to the product of two new couplings, for example $y_{L,11}y_{L,12}$. It is easy to check that this product is either zero or very small in all of our models. Nonetheless, it is interesting to estimate the size of such contributions in general. Clearly, for sufficiently high scales M, dimension-5 terms mediated by D-fermion exchange dominate over the dimension-6 D-scalar mediated terms, with the external scalars "dressed" by the appropriate superpartner loop. As explained in section 2, in the models we consider, the messenger scale M is sufficiently high such that the 2-loop Yukawa-mediated contributions to the scalar masses, which are of order $\alpha^2/(16\pi^2)F^2/M^2$ dominate over the 1-loop contributions $\sim \alpha/(4\pi)F^4/M^6$, where the factors of α denote the relevant couplings. Thus we have $F/M^2 \leq \sqrt{\alpha/(4\pi)}$. It is easy to check that for this choice, the contributions of higher-dimension terms mediated by D exchange are smaller than the contributions from (4.1) that we estimated above.

5 Conclusions

We presented models in which slepton masses are generated by messenger fields, through gauge and superpotential interactions. If such spectra are measured at the LHC, the GMSB structure will be apparent in the gaugino spectrum, with the slepton masses clearly indicating some flavor-dependent mediation of supersymmetry breaking, and providing additional handles on the source of fermion masses in the standard model. We concentrated on slepton masses, but as we explained, it is straightforward to generalize this construction to include messenger couplings to squarks.

It would also be interesting to examine mechanisms for generating the mu term in these models, since the flavor symmetries we discussed often forbid this term. The models may also accommodate large couplings of the Higgs to the supersymmetry-breaking sector in the spirit of [40].

Finally, while our models are based on flavor symmetries, it would be interesting to consider alternative frameworks for controlling both the Yukawa couplings and the matter-messenger couplings.

JHEP06(2012)124

Acknowledgments

We thank Yossi Nir, Yuval Grossman, and Yuri Shirman for helpful discussions. Parts of this work were done during a sabbatical year of YS at UCI. YS thanks the UCI hep group for its warm hospitality and financial support during this year. The work of YS and PZS was supported by the Israel Science Foundation (ISF) under Grant No. 1155/07, and by the United States-Israel Binational Science Foundation (BSF) under Grant No. 2006071.

References

[1] MEGA collaboration, M. Brooks et al., *New limit for the family number nonconserving decay* $\mu^+ \to e^+\gamma$, *Phys. Rev. Lett.* **83** (1999) 1521 [hep-ex/9905013] [INSPIRE].

[2] BABAR collaboration, B. Aubert et al., *Searches for lepton flavor violation in the decays* $\tau^{+-} \to e^{+-}\gamma$ *and* $\tau^{+-} \to \mu^{+-}\gamma$, *Phys. Rev. Lett.* **104** (2010) 021802 [arXiv:0908.2381] [INSPIRE].

[3] J.L. Feng, C.G. Lester, Y. Nir and Y. Shadmi, *The standard model and supersymmetric flavor puzzles at the large hadron collider*, *Phys. Rev.* **D 77** (2008) 076002 [arXiv:0712.0674] [INSPIRE].

[4] G.D. Kribs, E. Poppitz and N. Weiner, *Flavor in supersymmetry with an extended R-symmetry*, *Phys. Rev.* **D 78** (2008) 055010 [arXiv:0712.2039] [INSPIRE].

[5] Y. Nomura and D. Stolarski, *Naturally flavorful supersymmetry at the LHC*, *Phys. Rev.* **D 78** (2008) 095011 [arXiv:0808.1380] [INSPIRE].

[6] Y. Nomura, M. Papucci and D. Stolarski, *Flavorful supersymmetry from higher dimensions*, *JHEP* **07** (2008) 055 [arXiv:0802.2582] [INSPIRE].

[7] C. Gross and G. Hiller, *Flavorful hybrid anomaly-gravity mediation*, *Phys. Rev.* **D 83** (2011) 095015 [arXiv:1101.5352] [INSPIRE].

[8] M. Dine, A.E. Nelson and Y. Shirman, *Low-energy dynamical supersymmetry breaking simplified*, *Phys. Rev.* **D 51** (1995) 1362 [hep-ph/9408384] [INSPIRE].

[9] M. Dine, A.E. Nelson, Y. Nir and Y. Shirman, *New tools for low-energy dynamical supersymmetry breaking*, *Phys. Rev.* **D 53** (1996) 2658 [hep-ph/9507378] [INSPIRE].

[10] G. Dvali, G. Giudice and A. Pomarol, *The* μ *problem in theories with gauge mediated supersymmetry breaking*, *Nucl. Phys.* **B 478** (1996) 31 [hep-ph/9603238] [INSPIRE].

[11] M. Dine, Y. Nir and Y. Shirman, *Variations on minimal gauge mediated supersymmetry breaking*, *Phys. Rev.* **D 55** (1997) 1501 [hep-ph/9607397] [INSPIRE].

[12] G. Giudice and R. Rattazzi, *Extracting supersymmetry breaking effects from wave function renormalization*, *Nucl. Phys.* **B 511** (1998) 25 [hep-ph/9706540] [INSPIRE].

[13] Z. Chacko and E. Ponton, *Yukawa deflected gauge mediation*, *Phys. Rev.* **D 66** (2002) 095004 [hep-ph/0112190] [INSPIRE].

[14] F. Joaquim and A. Rossi, *Gauge and Yukawa mediated supersymmetry breaking in the triplet seesaw scenario*, *Phys. Rev. Lett.* **97** (2006) 181801 [hep-ph/0604083] [INSPIRE].

[15] F. Joaquim and A. Rossi, *Phenomenology of the triplet seesaw mechanism with gauge and Yukawa mediation of SUSY breaking*, *Nucl. Phys.* **B 765** (2007) 71 [hep-ph/0607298] [INSPIRE].

JHEP06(2012)124

[16] A. Brignole, F.R. Joaquim and A. Rossi, *Beyond the standard seesaw: neutrino masses from Kähler operators and broken supersymmetry*, *JHEP* **08** (2010) 133 [arXiv:1007.1942] [INSPIRE].

[17] C. Froggatt and H.B. Nielsen, *Hierarchy of quark masses, Cabibbo angles and CP-violation*, *Nucl. Phys.* **B 147** (1979) 277 [INSPIRE].

[18] N. Arkani-Hamed, H.-C. Cheng, J.L. Feng and L.J. Hall, *Probing lepton flavor violation at future colliders*, *Phys. Rev. Lett.* **77** (1996) 1937 [hep-ph/9603431] [INSPIRE].

[19] N. Arkani-Hamed, J.L. Feng, L.J. Hall and H.-C. Cheng, *CP violation from slepton oscillations at the LHC and NLC*, *Nucl. Phys.* **B 505** (1997) 3 [hep-ph/9704205] [INSPIRE].

[20] K. Agashe and M. Graesser, *Signals of supersymmetric lepton flavor violation at the CERN LHC*, *Phys. Rev.* **D 61** (2000) 075008 [hep-ph/9904422] [INSPIRE].

[21] J. Hisano, R. Kitano and M.M. Nojiri, *Slepton oscillation at large hadron collider*, *Phys. Rev.* **D 65** (2002) 116002 [hep-ph/0202129] [INSPIRE].

[22] J. Hisano, M.M. Nojiri and W. Sreethawong, *Discriminating electroweak-ino parameter ordering at the LHC and its impact on LFV studies*, *JHEP* **06** (2009) 044 [arXiv:0812.4496] [INSPIRE].

[23] R. Kitano, *A clean slepton mixing signal at the LHC*, *JHEP* **03** (2008) 023 [arXiv:0801.3486] [INSPIRE].

[24] J.L. Feng, S.T. French, C.G. Lester, Y. Nir and Y. Shadmi, *The shifted peak: resolving nearly degenerate particles at the LHC*, *Phys. Rev.* **D 80** (2009) 114004 [arXiv:0906.4215] [INSPIRE].

[25] J.L. Feng, S.T. French, I. Galon, C.G. Lester, Y. Nir, et al., *Measuring slepton masses and mixings at the LHC*, *JHEP* **01** (2010) 047 [arXiv:0910.1618] [INSPIRE].

[26] J.L. Feng, I. Galon, D. Sanford, Y. Shadmi and F. Yu, *Three-body decays of sleptons with general flavor violation and left-right mixing*, *Phys. Rev.* **D 79** (2009) 116009 [arXiv:0904.1416] [INSPIRE].

[27] B. Allanach, J. Conlon and C. Lester, *Measuring smuon-selectron mass splitting at the CERN LHC and patterns of supersymmetry breaking*, *Phys. Rev.* **D 77** (2008) 076006 [arXiv:0801.3666] [INSPIRE].

[28] S. Kaneko, J. Sato, T. Shimomura, O. Vives and M. Yamanaka, *Measuring lepton flavour violation at LHC with long-lived slepton in the coannihilation region*, *Phys. Rev.* **D 78** (2008) 116013 [arXiv:0811.0703] [INSPIRE].

[29] A. De Simone, J. Fan, V. Sanz and W. Skiba, *Leptogenic supersymmetry*, *Phys. Rev.* **D 80** (2009) 035010 [arXiv:0903.5305] [INSPIRE].

[30] A.J. Buras, L. Calibbi and P. Paradisi, *Slepton mass-splittings as a signal of LFV at the LHC*, *JHEP* **06** (2010) 042 [arXiv:0912.1309] [INSPIRE].

[31] T. Ito, R. Kitano and T. Moroi, *Measurement of the superparticle mass spectrum in the long-lived Stau scenario at the LHC*, *JHEP* **04** (2010) 017 [arXiv:0910.5853] [INSPIRE].

[32] R. Fok and G.D. Kribs, *μ to e in R-symmetric supersymmetry*, *Phys. Rev.* **D 82** (2010) 035010 [arXiv:1004.0556] [INSPIRE].

[33] H. Dreiner, S. Grab and T. Stefaniak, *Discovery potential of selectron or smuon as the lightest supersymmetric particle at the LHC*, *Phys. Rev.* **D 84** (2011) 035023 [arXiv:1102.3189] [INSPIRE].

JHEP06(2012)124

[34] G. D'Ambrosio, G. Giudice, G. Isidori and A. Strumia, *Minimal flavor violation: an effective field theory approach*, *Nucl. Phys.* **B 645** (2002) 155 [hep-ph/0207036] [INSPIRE].

[35] Y. Nir and N. Seiberg, *Should squarks be degenerate?*, *Phys. Lett.* **B 309** (1993) 337 [hep-ph/9304307] [INSPIRE].

[36] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, *A complete analysis of FCNC and CP constraints in general SUSY extensions of the standard model*, *Nucl. Phys.* **B 477** (1996) 321 [hep-ph/9604387] [INSPIRE].

[37] Y. Nir, *Probing new physics with flavor physics (and probing flavor physics with new physics)*, arXiv:0708.1872 [INSPIRE].

[38] M. Ciuchini, A. Masiero, P. Paradisi, L. Silvestrini, S. Vempati, et al., *Soft SUSY breaking grand unification: leptons versus quarks on the flavor playground*, *Nucl. Phys.* **B 783** (2007) 112 [hep-ph/0702144] [INSPIRE].

[39] G. Hiller, Y. Hochberg and Y. Nir, *Flavor changing processes in supersymmetric models with hybrid gauge- and gravity-mediation*, *JHEP* **03** (2009) 115 [arXiv:0812.0511] [INSPIRE].

[40] C. Csáki, A. Falkowski, Y. Nomura and T. Volansky, *New approach to the $\mu - B_\mu$ problem of gauge-mediated supersymmetry breaking*, *Phys. Rev. Lett.* **102** (2009) 111801 [arXiv:0809.4492] [INSPIRE].

JHEP06(2012)124