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Lepton flavour violation in unified models with $U(1)$-family symmetries [1]

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Abstract

Lepton flavour non-conserving processes are examined in the context of unified models with $U(1)$-family symmetries which reproduce successfully the low-energy hierarchy of the fermion mass spectrum and the Kobayashi - Maskawa mixing. These models usually imply mixing effects in the supersymmetric scalar sector. We construct the fermion and scalar mass matrices in two viable models, and calculate the mixing effects on the $\mu \to e\gamma$, $\mu \to 3e$ and $\tau \to \mu\gamma$ rare decays. The relevant constraints on the sparticle mass spectrum as well as the role of various MSSM parameters are discussed. © 1998 Elsevier Science B.V. All rights reserved.

One of the most dramatic consequences of supersymmetric extensions of the Standard Model (SM) is the appearance of new sources of flavour violations [1–3]. Supersymmetric partners of gauge, fermion and scalar fields generate new types of flavour violating diagrams at the one-loop level, which enhance considerably the various rare processes.

Flavour non-conserving processes may still be relatively suppressed if the matrices of the supersymmetric partners of fermions, i.e. those of scalar quarks and scalar leptons, are diagonal in flavour space. It is widely believed however, that a realistic spectrum for the fermion mass matrices can be obtained when additional symmetries discriminate between the various families of the known fermion fields. Such symmetries imply also a non-trivial structure for the corresponding scalar mass matrices. Rare processes, being sensitive to these changes, usually lead to hard violation of flavour.

In this work, we compute the branching ratios for lepton flavour violating decays in realistic models whose fermion and scalar mass textures are obtained by $U(1)$-family ($U(1)_f$) symmetries. In our analysis we choose both, symmetric and non-symmetric fermion textures by appropriate selection of the $U(1)_f$ fermion charges. The choice of the lepton sector to test the predictability, and possibly the viability, of $U(1)_f$-models is ideal as there are no large uncertainties (unlike the case of the quark sector where large ambiguities enter due to poor knowledge of hadronic matrix elements) in the calculations. In a previous work [4], we have given some estimates for the $\mu \to e\gamma$ decay is the context of a simple $U(1)_f$-model in the case of small $\tan\beta$ regime. Here, we extend

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our previous analysis and examine cases for large and small $\tan\beta$ and various values of the gravitino and gaugino masses. We use two-loop analysis for gauge couplings and take into account threshold effects to calculate the sparticle spectrum used to construct the relevant scalar and fermion mass matrices entering the above processes. We find that the non-observation of lepton flavour violating processes put rather strong lower limits on the sparticle mass spectrum, in particular when $\tan\beta$ is large.

A wide class of models, which naturally incorporate flavour non-diagonal scalar mass matrices, arises in string scenarios where the usual gauge symmetry is accompanied by a number of $U(1)$ factors, the latter playing the role of family symmetries. The fermion mass textures of the above models are dictated by the particular charges of the particles under the $U(1)_f$ symmetries, the specific flat direction which has been chosen as well as the string selection rules and other string symmetries. In general, there are only few tree-level couplings in the superpotential (usually only those responsible for the top, bottom and tau masses), while all other fermion mass entries are supposed to be generated by higher non-renormalizable terms.

Once the flat directions and the $U(1)_f$ charges of a particular model have been fixed, the scalar mass matrix structure may also be easily computed through the Kähler function

$$\mathcal{G} = \mathcal{K} + \log|\mathcal{W}|^2 \tag{1}$$

where $\mathcal{W}$ is the superpotential and $\mathcal{K}$ has the general form

$$\mathcal{K} = -\log(S+\bar{S}) - \sum h_n \log\left(T_n + \bar{T}_n\right) + Z_{i\bar{j}}\left(T_n, \bar{T}_n\right) Q_i \bar{Q}_j + \cdots \tag{2}$$

with Q_i being the matter fields, S the dilaton, whereas T_n are the other moduli fields. The scalar mass matrices are determined by $Z_{i\bar{j}}$ and $\mathcal{W}$. The form of the $Z_{i\bar{j}}$ function is dictated by the modular symmetries and depends on the moduli and the modular weights of the fields. Thus, at the tree level, the diagonal terms are the only non-zero entries in the scalar mass matrices. Higher order terms allowed by the symmetries of the specific model fill in the non-diagonal entries. In what follows, we will explore the flavour violating processes in two different models which give realistic fermion mass spectrum. In particular, we will calculate the $\mu \to e\gamma$, $\tau \to \mu\gamma$ and $\mu \to 3e$ processes in the context of supersymmetric models whose low energy theory is the MSSM model augmented by a $U(1)$ family symmetry. One of them is using a charge assignment where symmetric mass matrices appear, while the other assumes $U(1)$-charges which lead to non-symmetric textures.

We start with some preliminary remarks about the sources of flavour violations and set our notation and conventions. After the breaking of some possible unified symmetry to that of the Standard Model, lepton flavour violating Yukawa interaction which appears in the superpotential is

$$\mathcal{W} = e^{cT}\lambda_e \ell H + \cdots \tag{3}$$

where ℓ is the left lepton doublet, e is the right singlet lepton, H is the higgs doublet and λ_e represents the Yukawa coupling matrix in flavour space. In addition, soft supersymmetry breaking terms generate mass matrices for the charged slepton fields, denoted by $\tilde{m}_\ell$, $\tilde{m}_{e_R}$. The Yukawa and soft scalar mass-squared matrices are diagonalized by unitary transformations

$$\lambda_e = V_R^* \lambda_e^\delta V_L^\dagger \tag{4}$$

$$\tilde{m}_\ell^2 = U_\ell\left(\tilde{m}_\ell^2\right)^\delta U_\ell^\dagger, \quad \text{and } \tilde{m}_{e_R}^2 = U_{e_R}\left(\tilde{m}_{e_R}^2\right)^\delta U_{e_R}^\dagger \tag{5}$$

where δ denotes diagonal. The lepton mass eigenstates (eig) are related to the weak eigenstates (w) by

$$\ell_{eig} = V_L^\dagger \ell_w, \quad e_{eig}^c = V_R^\dagger e_w^c \tag{6}$$

The charged slepton mass-squared matrix is a 6×6 matrix, built up from the two 3×3 left $\tilde{m}_\ell^2$ and right $\tilde{m}_{e_R}^2$ soft ones, as well as the off-diagonal submatrix which has the form

$$\left(m_{LR}^\ell\right)^2 = m_\ell(\mu\tan\beta + A_\ell) \tag{7}$$

where $m_\ell = \lambda_e \upsilon \cos\beta/\sqrt{2}$ is the charged lepton mass matrix, A_ℓ and μ are the trilinear and higgs mixing parameters in the superpotential, $\tan\beta$ the ratio of the two higgs vev's and $\upsilon = 246\text{GeV}$.

To calculate the mixing effects in the amplitudes, we work in the basis where the fermion mass matrix is diagonal,

$$\begin{aligned}(e^{cT})_w m_\ell \ell_w &= (e^{cT})_w V_R^* m_\ell^\delta V_L^\dagger \ell_w \\ &= (e^c)_{eig}^T m_\ell^\delta \ell_{eig}\end{aligned} \tag{8}$$

In this basis, the off-diagonal term (7) is written

$$m_{\ell}(\mu\tan\beta+A_{\ell}) = V_R^{*}\, m_{\ell}^{\delta}(\mu\tan\beta+A_{\ell})V_L^{\dagger} \qquad (9)$$

This defines uniquely the scalars in the basis where the fermions are in their mass-eigenstates

$$\tilde{e}_R^{*\prime} = \left(V_R^{\dagger}(\tilde{e}_R)\right)_w^{*} = V_R^{T}(\tilde{e}_R)_w^{*} \qquad (10)$$

$$\tilde{\ell}' = V_L^{\dagger}\tilde{\ell}_w. \qquad (11)$$

The soft terms for right and left charged sleptons must be written in the same basis. The 6×6 matrix then takes the form

$$\begin{pmatrix} V_L^{\dagger}\tilde{m}_{\ell}^2 V_L & \left((A_{\ell}+\mu\tan\beta)m_{\ell}^{\delta}\right)^{\dagger} \\ (A_{\ell}+\mu\tan\beta)m_{\ell}^{\delta} & V_R^{T}\tilde{m}_{e_R}^2 V_R^{*} \end{pmatrix}. \qquad (12)$$

A well known result in the context of the non-supersymmetric standard model is the conservation of lepton flavour in the case of zero neutrino masses, while in the case of massive, non-degenerate neutrinos, the amount of lepton flavour violation is proportional to the factor $\Delta m_{\nu}^2/M_W^2$ [5], which highly suppresses all relevant processes. When supersymmetry enters the game, the whole scene changes completely. Even in the absence of right handed neutrinos, flavour violations could occur via the exchange of supersymmetric particles. A large number of new parameters (sparticle masses, mixing angles, e.t.c.) appear in the calculations, enlarging therefore the parameter space and making difficult the consistency of the predicted branching ratios with the experimental bounds. In the context of unification and low energy phenomenology scenarios, these processes can provide useful constraints on the parameter space.

We will briefly present the minimum number of inputs necessary to determine all low energy parameters entering in a lepton flavour violating process. In the context of supersymmetric unified models, we assume a universality condition for the scalar masses at the unification scale M_U. The general formula, at this scale, is $\tilde{m}_i^2(M_U)=(1+q_i)m_{3/2}^2$, where $m_{3/2}$ is the gravitino mass, i is a flavour index and q_i is the modular weight of the corresponding field. This tree-level contribution is flavour diagonal. Non-diagonal terms are expected to appear through non-renormalizable terms with the expense of an extra parameter ϵ, namely $\epsilon=\langle\phi\rangle/M$ where $\langle\phi\rangle$ is a singlet field vev and M a Planck-scale mass. The magnitude of the vev $\langle\phi\rangle$ can be fit from the fermion sector. A crucial role is also played by the gaugino soft masses, the trilinear soft parameter A as well as the Yukawa couplings λ_t and λ_b at the unification scale M_U (we assume that $\lambda_b(M_U)=\lambda_{\tau}(M_U)$). In the minimal scenario, the gaugino masses at the unification scale are determined in terms of the universal mass parameter $m_{1/2}$. Thus, at M_U, we use a minimum set of parameters, namely $(m_{1/2},m_{3/2},\mu,A,\lambda_t,\lambda_b)$, together with the value of the common coupling α_U and the unification scale M_U in such a way that after the renormalization group running we obtain a consistent set of all low energy measured quantities. For any acceptable such set, we calculate the branching ratios of the flavour violating processes.

Fig. 1 shows the one-loop diagrams relevant to the $\mu\to e\gamma$ process. The corresponding $\tau\to\mu\gamma$-decay is represented by an analogous set of graphs. $\mu\to 3e$ proceeds through the decay of the (now virtual) photon to an electron-positron pair. There are also box-diagrams contributing to this process, they are however relatively suppressed.

The electromagnetic current operator between two lepton states l_i and l_j is given in general by

$$\begin{aligned}\mathcal{T}_{\lambda} &= \langle l_i|(p-q)|\mathcal{J}_{\lambda}|l_j(p)\rangle \\ &= \bar{u}_i(p-q)\Big\{m_j i\sigma_{\lambda\beta}q^{\beta}\left(A_M^L P_L + A_M^R P_R\right) \\ &\quad + \left(q^2\gamma_{\lambda}-q_{\lambda}\gamma\cdot q\right)\left(A_E^L P_L + A_E^R P_R\right)\Big\}u_j(p)\end{aligned} \qquad (13)$$

where q is the photon momentum. The A_M's and A_E's have contributions from neutralino-charged slepton (n) and chargino-sneutrino (c) exchange

$$A_M^{L,R} = A_{M(n)}^{L,R} + A_{M(c)}^{L,R}, \quad A_E^{L,R} = A_{E(n)}^{L,R} + A_{E(c)}^{L,R} \qquad (14)$$

The amplitude of the process is then proportional to $\mathcal{T}_{\lambda}\epsilon^{\lambda}$ where ϵ^{λ} is the photon polarization vector. The easiest way to determine the loop momentum integral contribution to the A's is to search, in the corresponding diagram, for terms proportional to $(p\cdot\epsilon)$ and $(q\cdot\epsilon)$. The coefficient of the former is

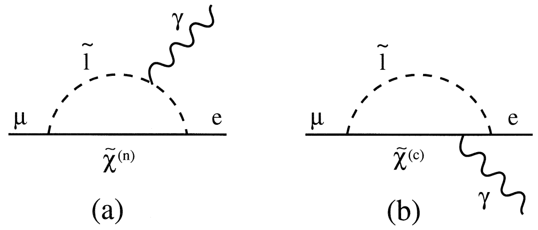

Fig. 1. The generic Feynman diagrams for the $\mu \to e\gamma$ decay. $\tilde{l}$ stands for charged slepton (a) or sneutrino (b), while $\tilde{\chi}^{(n)}$ and $\tilde{\chi}^{(c)}$ represent neutralinos and charginos respectively.

proportional to the momentum integral contribution to the $\sigma_{\lambda\beta}$ term in (13), while the coefficient of the latter is proportional to the difference of the momentum integral contribution between the $\sigma_{\lambda\beta}$ and the $(q^2\gamma_\lambda - q_\lambda \gamma\cdot q)$ terms. Defining the ratio $x = M^2/m^2$, where M is the chargino (neutralino) mass and m the sneutrino (charged slepton) mass, the following functions appear in the A_M term

$$A_{M(n)}: \quad \frac{1}{6(1-x)^4}\left(1-6x+3x^2+2x^3 - 6x^2\log x\right)$$

and

$$\frac{1}{(1-x)^3}\left(1-x^2+2x\log x\right)\frac{M}{m_{l_j}}$$

$$A_{M(c)}: \quad \frac{1}{6(1-x)^4}\left(2+3x-6x^2+x^3+6x\log x\right)$$

and

$$\frac{1}{(1-x)^3}\left(-3+4x-x^2-2\log x\right)\frac{M}{m_{l_j}} \qquad (15)$$

where m_{l_j} is the mass of the l_j lepton, while for the A_E we have

$$A_{E(n)}: \quad \frac{1}{(1-x)^4}\left(2-9x+18x^2-18x^3 + 6x^3\log x\right)$$

$$A_{E(c)}: \quad \frac{1}{(1-x)^4}\left(16-45x+36x^2-7x^3 + 6(2-3x)\log x\right) \qquad (16)$$

Notice the lack of terms proportional to the gaugino mass M which cancel. The Branching Ratio (BR) of the decay $l_j \to l_i + \gamma$ is given by

$$BR(l_j \to l_i\gamma) = \frac{48\pi^3\alpha}{G_F^2}\left(\left(A_M^L\right)^2 + \left(A_M^R\right)^2\right)$$

Our approach to determine the mixing effects is the following:

- For each particular model we construct the lepton and left and right slepton mass matrices and determine the corresponding diagonalizing matrices.
- From the input parameters at M_U (α_U, M_U, $m_{1/2}$, e.t.c.) and using the RGEs, we determine the soft masses for gauginos and sleptons, the Yukawa couplings λ_t and λ_b and the μ and A parameters. We are using two-loop-β functions and incorporate threshold effects for the scalars and gauginos. In Table 1 we show the results of the RGEs running for four characteristic cases (In the runnings we have assumed the condition $m_{1/2} > m_{3/2}$. We have not explored the parameter space where the inverse relation holds).
- Using the above values and the form of the mass matrices for gauginos and sleptons we determine the corresponding mass eigenvalues and eigenstates. The diagonalizing matrices determined in the first step are used to transform the slepton mass matrices in the desired basis where the lepton mass matrix is diagonal.
- Having all relevant diagonalizing matrices (for charginos, neutralinos and sleptons), plus the eigenvalues of the gauginos and sleptons we can construct the amplitude of the process.

We present here the relevant mass matrices of two models whose successful fermion mass hierarchy is predicted by $U(1)$ symmetries. As a first example, we use the scalar mass matrices obtained in a simple $SU(3)\times SU(2)\times U(1)$ model with an additional $U(1)_f$ symmetry [6]. After the implementation of this symmetry, the fermion matrix for leptons in this model is given

$$m_\ell \approx \begin{pmatrix} \tilde{\epsilon}^{2|a+b|} & \tilde{\epsilon}^{|a|} & \tilde{\epsilon}^{|a+b|} \\ \tilde{\epsilon}^{|a|} & \tilde{\epsilon}^{2|b|} & \tilde{\epsilon}^{|b|} \\ \tilde{\epsilon}^{|a+b|} & \tilde{\epsilon}^{|b|} & 1 \end{pmatrix} m_\tau \qquad (17)$$

Table 1
Inputs and outputs of the RGEs running for four representative cases (masses are in GeV)

	Case			
	(a)	(b)	(c)	(d)
Inputs				
α_U^{-1}, $M_U/10^{16}$	24.88, 1.0	25.07, 0.67	25.00, 0.67	25.00, 0.67
$m_{3/2}$, $m_{1/2}$	138, 350	150, 380	160, 390	160, 390
λ_t, λ_b	0.85, 0.10	0.90, 0.10	1.05, 0.08	1.05, 0.05
A, μ	−207, −395	−225, −395	−240, −395	−240, −395
Results				
M_1, M_2, M_3	151, 295, 820	166, 321, 865	171, 329, 888	170, 329, 890
$\tilde m^L_{(1,2)}$, $\tilde m^L_{(3)}$	281, 279	303, 301	314, 313	316, 314
$\tilde m^l_{(1,2)}$, $\tilde m^l_{(3)}$	193, 187	209, 203	219, 215	219, 218
$\tilde m^Q_{1,2}$, $\tilde m^Q_3$	750, 678	789, 712	811, 730	813, 733
$\tilde m^u_{1,2}$, $\tilde m^u_3$	718, 564	754, 589	775, 598	778, 599
$\tilde m^d_{1,2}$, $\tilde m^d_3$	714, 199	749, 219	771, 244	773, 252
A_τ, μ	10.3, −324	8.3, −319	3.3, −299	6.9, −301
$\tan\beta$, m_t	14, 185	14, 188	11, 195	7, 194

where the parameter $\tilde\epsilon$ is some power of the singlet vev scaled by the unification mass, while a,b are certain combinations of the lepton and quark $U(1)_f$-charges. Order one parameters in front of the various entries (not calculable in this simple model) are assumed to reproduce exactly the fermion mass relations after renormalization group running. [2] The scalar mass matrices of this model are given in [4]. For the sleptons we obtain

$$\tilde m^2_{\ell, e_R} \approx \begin{pmatrix} 1 & \tilde\epsilon^{\,|a+6b|} & \tilde\epsilon^{\,|a+b|} \\ \tilde\epsilon^{\,|a+6b|} & 1 & \tilde\epsilon^{\,|b|} \\ \tilde\epsilon^{\,|a+b|} & \tilde\epsilon^{\,|b|} & 1 \end{pmatrix} m^2_{3/2} \qquad (18)$$

A successful lepton mass hierarchy in this case is obtained for the choice $a=3, b=1$ and $\tilde\epsilon = 0.23$.

As a second example in the calculation of the rare processes, we apply our results in several textures obtained in context of the $SU(4)$ model in Ref. [7]. As far as the symmetric fermion mass textures are concerned, the analysis follows the same lines as in the first example, since the matrices obtained are similar to those above. The non-symmetric case is however completely different. There, the violations are much harder and the bounds on the various scalar mass parameters increase substantially. In order to be specific, we work out a particular example based on the charge assignment of Table 2. The field assignment is $F+\bar F = (4,2,1) + (\bar 4,1,2)$ for the three generations, $H+\bar H(4,1,2) + (\bar 4,1,2)$ for the higgses and $h=(1,2,2)$ for the bidoublet including the standard higgs doublets. Notice that this $U(1)$ symmetry is anomalous. However the mixed anomalies are zero. This fact allows for a Green-Schwarz anomaly cancellation mechanism. The fermion mass matrices are generated by operators of the form

$$\bar F_i F_j h \delta^m \epsilon^n \qquad (19)$$

where $\delta = H\bar H$ is an effective singlet generated by the higgs fields $H, \bar H$.

Taking into account Clebsch-Gordan coefficients derived from these operators, as in Ref. [7] and using

Table 2
$U(1)_f$ charges of fields in the model

F_1	F_2	F_3	$\bar F_1$	$\bar F_2$	$\bar F_3$	h	H	$\bar H$
4	0	−1	3	−1	1	0	1	−1

[2] We use the Yukawa coefficients $c_{11}=4.0; c_{12}=c_{21}=0.9; c_{22}=1.08; c_{33}=1.9$.

the values $\epsilon \sim 0.14$, $\delta \sim 0.21$, we arrive at the following lepton Yukawa matrix at the unification scale

$$\lambda_e = \begin{pmatrix} 0 & \sqrt{2}\,\eta_{12} & 0 \\ \sqrt{2}\,\eta_{21} & 3\eta_{22}/\sqrt{5} & 0 \\ 0 & -3\sqrt{2}\,\eta_{32}/\sqrt{5} & \eta_{33} \end{pmatrix} \tag{20}$$

The following unification scale input parameters give the correct lepton mass spectrum

$$\eta_{22} = 2.88 \times 10^{-2}, \quad \eta_{12} = 2.81 \times 10^{-3},$$
$$\eta_{21} = 1.30 \times 10^{-3}, \quad \eta_{33} = 1.18,$$
$$\eta_{32} = 7.28 \times 10^{-2}. \tag{21}$$

The scalar mass matrices are determined from the $U(1)$ charges chosen to give correct predictions for the fermions. Since left and right fields have different charges, we obtain two different types of scalar mass matrices

$$\tilde{m}_l^2 = \begin{pmatrix} 1 & \bar{\epsilon}^4 & \bar{\epsilon}^5 \\ \epsilon^4 & 1 & \bar{\epsilon} \\ \epsilon^5 & \epsilon & 1 \end{pmatrix} \tag{22}$$

$$\tilde{m}_{e_R}^2 = \begin{pmatrix} 1 & \bar{\epsilon}^4 & \bar{\epsilon}^2 \\ \epsilon^4 & 1 & \epsilon^2 \\ \epsilon^2 & \bar{\epsilon}^2 & 1 \end{pmatrix} \tag{23}$$

The two expansion parameters ϵ and $\bar{\epsilon}$ are defined as follows

$$\epsilon \equiv \langle \theta \rangle / M_U \sim \bar{\epsilon} = \langle \bar{\theta} \rangle / M_U \tag{24}$$

In Table 3 we show the branching ratios for the four cases appearing in Table 1 and for the two models discussed above. We also give the L and R amplitudes for the neutralino and chargino exchanges for each case (see Eqs. (13), (14)).

From Tables 1 and 3, we can conclude that lepton flavour violation in a viable class of supersymmetric unified theories puts non-trivial constraints on the scalar mass spectrum. In the case of symmetric textures of mass matrices, the $\mu \to e\gamma$-branching ratio is found to exceed the present experimental bound (4.9×10^{-11}) for values of the gravitino mass parameter $m_{3/2}$ less than about 150GeV, in particular when $\tan\beta$ obtains intermediate or higher values. In our approach, the gaugino mass parameter $m_{1/2}$ is also found to have a lower bound (≥ 350GeV) for consistency with the experimental $\mu \to e\gamma$-bound. In the case of non-symmetric masses matrices, the bounds are even higher: the three cases (a,b,c) of Table 3 violate the present experimental bounds, implying therefore a rather heavy supersymmetric mass spectrum. For small $\tan\beta$, however, we obtain results consistent with the experimental bounds as in case (d). Actually, we see a strong dependence of the branching ratio on $\tan\beta$. As $\tan\beta \to 1$, the neutralino exchange processes get smaller and smaller. The chargino ones, while individually remain of the same order, exhibit an increasing cancellation, rendering the branching ratio much lower that the experimental bound. The $\tau \to \mu\gamma$-rare decay (branching ratio $< 4.2 \times 10^{-6}$), does not put further constraints as can been checked from the same Table . Moreover, the branching ratio for the $\mu \to 3e$-decay in all cases is found about three orders smaller than $BR(\mu \to e\gamma)$.

Returning to the results of Table 1, we infer that non-observation of the $\mu \to e\gamma$-decay implies that all scalars appear with masses at least heavier than about 200GeV. Although our results are only for two specific examples, their main characteristics are rather generic and all sparticle mass bounds from flavour decays are expected always much larger than those obtained from other types of experiments. In this selected region of the parameter space (i.e. $m_{1/2} > m_{3/2}$) the sleptons are the lightest scalars. Clearly, the relatively large flavour violations are due the large $\tan\beta$-effects as well as to the fact that $U(1)_f$-models imply also mixing effects to the scalar sector. The rather high scalar mass bounds could be considered as an ominous perspective for models with large $\tan\beta$ and non-symmetric fermion mass textures, in particular by those who envisage a relatively light supersymmetric mass spectrum accessible to future experiments. We should note however, that slepton masses of this order are in the range of the LHC. Indeed, slepton decays can in principle be detected in CMS with a mass up to 350GeV [8] Finally, we wish to comment that in more complicated structures with cyclic permutation symmetries between generations and universal anomalous $U(1)$-factors may prevent mixing effects in the supersymmetric mass matrices [9]. In such cases, the above constraints are relaxed.

In conclusion, we have examined lepton flavour rare processes in a class of supersymmetric gauge

Table 3
Branching ratios for the two processes $\mu \to e\gamma$ and $\tau \to \mu\gamma$, corresponding to the four cases of Table 1 and for the two models with symmetric and non-symmetric mass textures. We also show the L and R amplitudes with neutralino and chargino exchanges

	Case			
	(a)	(b)	(c)	(d)
Symmetric textures in $U(1)_f$-models				
$A^L_{M(n)}$	$3.58 \cdot 10^{-11}$	$2.74 \cdot 10^{-11}$	$1.22 \cdot 10^{-11}$	$1.83 \cdot 10^{-12}$
$A^R_{M(n)}$	$-9.47 \cdot 10^{-12}$	$-6.79 \cdot 10^{-12}$	$-2.04 \cdot 10^{-12}$	$-1.11 \cdot 10^{-12}$
$A^L_{M(c)}$	$4.13 \cdot 10^{-14}$	$3.57 \cdot 10^{-14}$	$1.44 \cdot 10^{-14}$	$8.84 \cdot 10^{-15}$
$A^R_{M(c)}$	$8.43 \cdot 10^{-12}$	$7.30 \cdot 10^{-12}$	$2.93 \cdot 10^{-12}$	$1.79 \cdot 10^{-12}$
$BR(\mu \to e\gamma)$	$1.03 \cdot 10^{-10}$	$6.01 \cdot 10^{-11}$	$1.20 \cdot 10^{-11}$	$3.07 \cdot 10^{-13}$
$A^L_{M(n)}$	$9.15 \cdot 10^{-10}$	$7.05 \cdot 10^{-10}$	$3.23 \cdot 10^{-10}$	$4.98 \cdot 10^{-11}$
$A^R_{M(n)}$	$2.45 \cdot 10^{-10}$	$1.83 \cdot 10^{-10}$	$6.71 \cdot 10^{-11}$	$3.18 \cdot 10^{-11}$
$A^L_{M(c)}$	$-1.36 \cdot 10^{-11}$	$-1.18 \cdot 10^{-11}$	$-4.74 \cdot 10^{-12}$	$-2.91 \cdot 10^{-12}$
$A^R_{M(c)}$	$-2.26 \cdot 10^{-10}$	$-1.96 \cdot 10^{-10}$	$-7.87 \cdot 10^{-11}$	$-4.80 \cdot 10^{-11}$
$BR(\tau \to \mu\gamma)$	$6.50 \cdot 10^{-8}$	$3.84 \cdot 10^{-8}$	$8.08 \cdot 10^{-9}$	$1.96 \cdot 10^{-10}$
Non-symmetric fermion mass textures				
$A^L_{M(n)}$	$-4.55 \cdot 10^{-10}$	$-3.56 \cdot 10^{-10}$	$-1.70 \cdot 10^{-10}$	$-2.55 \cdot 10^{-11}$
$A^R_{M(n)}$	$6.98 \cdot 10^{-12}$	$5.30 \cdot 10^{-12}$	$1.72 \cdot 10^{-12}$	$1.59 \cdot 10^{-12}$
$A^L_{M(c)}$	$-2.41 \cdot 10^{-14}$	$-5.98 \cdot 10^{-14}$	$-1.22 \cdot 10^{-14}$	$-1.48 \cdot 10^{-12}$
$A^R_{M(c)}$	$-4.87 \cdot 10^{-12}$	$-1.21 \cdot 10^{-11}$	$-2.43 \cdot 10^{-12}$	$-2.97 \cdot 10^{-12}$
$BR(\mu \to e\gamma)$	$1.66 \cdot 10^{-8}$	$1.01 \cdot 10^{-9}$	$2.30 \cdot 10^{-9}$	$5.21 \cdot 10^{-11}$
$A^L_{M(n)}$	$-8.51 \cdot 10^{-10}$	$-6.56 \cdot 10^{-10}$	$-3.03 \cdot 10^{-10}$	$-4.39 \cdot 10^{-11}$
$A^R_{M(n)}$	$1.59 \cdot 10^{-10}$	$1.16 \cdot 10^{-10}$	$3.38 \cdot 10^{-11}$	$2.93 \cdot 10^{-11}$
$A^L_{M(c)}$	$-1.53 \cdot 10^{-11}$	$-1.32 \cdot 10^{-11}$	$-5.31 \cdot 10^{-12}$	$-3.27 \cdot 10^{-12}$
$A^R_{M(c)}$	$-2.55 \cdot 10^{-10}$	$-2.21 \cdot 10^{-10}$	$-8.86 \cdot 10^{-11}$	$-5.40 \cdot 10^{-11}$
$BR(\tau \to \mu\gamma)$	$6.08 \cdot 10^{-8}$	$3.67 \cdot 10^{-8}$	$7.83 \cdot 10^{-9}$	$2.27 \cdot 10^{-10}$

theories where an additional $U(1)$ symmetry discriminates the three families. Such symmetries, are capable of generating successfully the hierarchical fermion mass spectrum and the Kobayashi-Maskawa mixing in the hadronic sector, when non-renormalizable contributions are taken into account in the superpotential. They, however, imply mixing effects through non-renormalizable terms in the Kähler potential, and consequently, in the scalar partners of quarks and leptons, leading thus to hard flavour violations. These violations are even larger when models with non-symmetric textures and large $\tan\beta$ values are considered. As a result, stringent constraints are found in the sparticle spectrum. Such bounds could be relaxed if, for example, additional symmetries of the Kähler potential are imposed, so that the scalar partners of fermions remain flavour diagonal, while at the same time they do not appear in the superpotential. We hope to come back to this issue in a future publication.

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