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Minimal Flavour Violation and Neutrino Masses without R-parity

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Abstract

We study the extension of the Minimal Flavour Violation (MFV) hypothesis to the MSSM without R-parity. The novelty of our approach lies in the observation that supersymmetry enhances the global symmetry of the kinetic term and in the fact that we consider as irreducible sources of the flavour symmetry breaking all the couplings of the superpotential including the R-parity violating ones. If R-parity violation is responsible for neutrino masses, our setup can be seen as an extension of MFV to the lepton sector. We analyze two patterns based on the non-abelian flavour symmetries $SU(3)^4 \otimes SU(4)$ and $SU(3)^5$. In the former case the total lepton number and the lepton flavour number are broken together, while in the latter the lepton number can be broken independently by an abelian spurion, so that visible effects and peculiar correlations can be envisaged in flavour changing charged lepton decays like $\ell_i \to \ell_j \gamma$.

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1 Introduction

The flavour problem can be viewed as the clash between the theoretical expectation of New Physics (NP) at the TeV scale and the experimental observations in Flavour Changing Neutral Current (FCNC) processes which severely constrain the scale $\Lambda_{\rm NP}$ of the NP beyond the 10^4 TeV domain (for a review see e.g. Ref. [1]). If we insist in keeping $\Lambda_{\rm NP} \approx$ TeV for naturalness, then we have to conclude that the flavour structure of the NP is highly non-generic.

The Minimal Flavour Violation (MFV) hypothesis [2] is a powerful organizing principle which states that the sources of flavour symmetry breaking of the NP are aligned to the Standard Model (SM) Yukawas. This ansatz provides an automatic suppression of the NP contribution to the flavour violating observables and thus a solution of the aforementioned flavour problem (see for instance Ref. [3]).

If MFV is at play in the quark sector, it is reasonable then to assume it also for leptons. However, the extension of MFV to the lepton sector is less straightforward, since the mechanism itself generating neutrino masses is unknown and several scenarios can be envisaged. Starting from Ref. [4] many formulations of Minimal Lepton Flavour Violation (MLFV) have been proposed and analyzed [5, 6, 7, 8, 9, 10, 11].

In this work we consider another interesting possibility in the context of the Minimal Supersymmetric SM (MSSM) without R-parity (for a review see e.g. Ref. [12]). Our analysis is moved by two simple observations about the MSSM:

1. The largest group of unitary transformations commuting with the gauge group (and with supersymmetry) is $U(3)_{\hat{q}} \otimes U(3)_{\hat{u}^c} \otimes U(3)_{\hat{d}^c} \otimes U(3)_{\hat{e}^c} \otimes U(4)_{\hat{L}} \otimes U(1)_{\hat{h}_u}$.

   The presence of the $U(4)_{\hat{L}}$ factor is due to the fact that the superfields $\hat{\ell}$ and $\hat{h}_d$ have the same quantum numbers, so that it is possible to rearrange them into a 4-dimensional flavour multiplet $\hat{L}$.

2. The MSSM has already all the degrees of freedom sufficient to generate neutrino masses and mixings through R-parity Violating (RPV) interactions [13], without the need of any extra state.

Thus the aim of our work is twofold: we first generalize the MFV expansion of the soft terms by including also the RPV couplings as the original sources of flavor breaking and then we connect the RPV spurions with the neutrino sector observables, providing an alternative scenario of MLFV.

Our approach towards the R-parity differs from that of Refs. [14, 15] in the fact that we do not aim at an explanation of the smallness of the RPV couplings. We simply treat them, in a more democratic way, on the same ground of all the other couplings of the superpotential. Remarkably, the values of the RPV couplings needed in order to fit neutrino masses are of the same order of magnitude of the SM Yukawas of the first and second families.

In the following we analyze two symmetry patterns based on the flavour symmetries $SU(3)^4 \otimes SU(4)$ and $SU(3)^5 \otimes U(1)_L \otimes U(1)_B$. In order to exemplify the connection of the RPV spurions with the neutrino observables we introduce a toy model in which only μ^i and λ'^{i33} are switched on, though a similar analysis could be performed also in more realistic RPV models of neutrino masses.

In the case of the former flavour symmetry the breaking scale of lepton number is linked to that of lepton flavour violation (LFV), thus implying small effects in LFV physics. On the other hand the latter flavour symmetry allows to break the lepton number independently by means of an abelian spurion, so that visible effects are in principle achievable. We finally study the correlations among the flavour changing charged lepton decays $\ell_i \to \ell_j \gamma$.

2 Minimal Flavour Violation without R-parity

The starting point of the MFV idea is based on the observation that the largest group of unitary transformations commuting with the SM gauge group is

$$G^{\rm SM}_{\rm kin} = U(3)_q \otimes U(3)_{u^c} \otimes U(3)_{d^c} \otimes U(3)_{e^c} \otimes U(3)_\ell \otimes U(1)_h \, . \tag{1}$$

This corresponds to the global symmetry of the gauge invariant kinetic term of the SM fields

$$\Phi = \left(q_i, u^c_i, d^c_i, e^c_i, \ell_i, h\right) \, , \tag{2}$$

with i spanning over the three families. Notice that ℓ_i and h have the same gauge quantum numbers and only the Lorentz structure prevents the global symmetry of the kinetic term from being larger.

On the other hand the situation in the MSSM is qualitatively different since the supersymmetrization of the SM spectrum restore the symmetry between scalars and fermions, thus enhancing the global symmetry of the kinetic term.

In order to make apparent this enhancement it is useful to define a generalized lepton multiplet $\hat{L}_\alpha = (\hat{\ell}_i, \hat{h}_d)$ and rewrite the set of chiral superfields of the MSSM in the following way

$$\hat{\Phi} = \left(\hat{q}_i, \hat{u}^c_i, \hat{d}^c_i, \hat{e}^c_i, \hat{L}_\alpha, \hat{h}_u\right) \, , \tag{3}$$

where a second Higgs doublet is introduced in order to ensure anomaly cancellation. Then the global symmetry of the kinetic term

$$\int d^4\theta \; \hat{\Phi}^\dagger e^{2g\hat{V}} \hat{\Phi} \tag{4}$$

turns out to be

$$G_{\rm kin}^{\rm MSSM} = U(3)_{\hat q} \otimes U(3)_{\hat u^c} \otimes U(3)_{\hat d^c} \otimes U(3)_{\hat e^c} \otimes U(4)_{\hat L} \otimes U(1)_{\hat h_u} . \tag{5}$$

Notice that this holds irrespectively of the fact that R-parity is or not an exact symmetry of the full MSSM lagrangian.

After decomposing $G_{\rm kin}^{\rm MSSM}$ in abelian and non-abelian factors we can identify a linear combination of the six $U(1)$ generators with the generator of SM hypercharge. It is possible then to define the generalized flavour group of the MSSM as

$$G_F = SU(3)_{\hat q} \otimes SU(3)_{\hat u^c} \otimes SU(3)_{\hat d^c} \otimes SU(3)_{\hat e^c} \otimes SU(4)_{\hat L} , \tag{6}$$

while the abelian factors can be rearranged in the following way

$$G_A = U(1)_{\hat u^c} \otimes U(1)_{\hat d^c} \otimes U(1)_{\hat e^c} \otimes U(1)_{\hat L} \otimes U(1)_B , \tag{7}$$

where B is the baryon number. G_F and G_A are explicitly broken by the most general MSSM superpotential and soft lagrangian.

Since the MSSM has many sources of flavour violation it is useful to have a rationale in order to select the origin of this breaking. Let us imagine that the flavour symmetry is broken at the scale Λ_F by some unknown mechanism. Then, if the breaking of SUSY is due to a flavour universal mechanism (like in gauge mediation [16]) and the scale of mediation M is smaller than Λ_F, it is natural to expect that the soft terms feel the breaking of flavour only through supersymmetric interactions.

Having in mind such a MFV framework we *assume* that the original source of flavour violation is given by the the couplings of the most general MSSM superpotential

$$W = Y_U^{ij} \hat q_i \hat u^c_j \hat h_u + Y_D^{\alpha ij} \hat L_\alpha \hat q_i \hat d^c_j + \tfrac{1}{2} Y_E^{\alpha\beta i} \hat L_\alpha \hat L_\beta \hat e^c_i + \mu^\alpha \hat h_u \hat L_\alpha + \tfrac{1}{2} (\lambda'')^{ijk} \hat u^c_i \hat d^c_j \hat d^c_k , \tag{8}$$

where the gauge structure has been omitted for simplicity. Notice also the antisymmetry of the couplings $Y_E^{\alpha\beta i} = -Y_E^{\beta\alpha i}$ and $(\lambda'')^{ijk} = -(\lambda'')^{ikj}$.

In order to formally restore the invariance with respect to the flavor group we treat the couplings in Eq. (8) as spurions, with quantum numbers under G_F:

$$\begin{aligned}
Y_U &\sim (\bar 3, \bar 3, 1, 1, 1) && (9)\\
Y_D &\sim (\bar 3, 1, \bar 3, 1, \bar 4) && (10)\\
Y_E &\sim (1, 1, 1, \bar 3, 6) && (11)\\
\mu &\sim (1, 1, 1, 1, \bar 4) && (12)\\
\lambda'' &\sim (1, \bar 3, 3, 1, 1) , && (13)
\end{aligned}$$

where our conventions are such that each chiral superfield in $\hat\Phi$ transforms according to the fundamental representation of the relative group factor of G_F.

Following the MFV principle we can expand the soft terms (cf. Appendix A for the notation) by means of the spurions in Eqs. (9)–(13)

$$
\begin{aligned}
\left(\tilde{m}^2_q\right)^i_j &= \tilde{m}^2\left(c_q\delta^i_j + d^1_q\ Y_U^{ik}\left(Y_U^*\right)_{jk} + d^2_q\ Y_D^{\alpha ik}\left(Y_D^*\right)_{\alpha jk}\right)\\
\left(\tilde{m}^2_{u^c}\right)^i_j &= \tilde{m}^2\left(c_{u^c}\delta^i_j + d^1_{u^c}\ Y_U^{ki}\left(Y_U^*\right)_{kj} + d^2_{u^c}(\lambda'')^{ikl}(\lambda''^*)_{jkl}\right)\\
\left(\tilde{m}^2_{d^c}\right)^i_j &= \tilde{m}^2\left(c_{d^c}\delta^i_j + d^2_{d^c}\ Y_D^{\alpha ki}\left(Y_D^*\right)_{\alpha kj} + d^2_{d^c}(\lambda'')^{kil}(\lambda''^*)_{kjl}\right)\\
\left(\tilde{m}^2_{e^c}\right)^i_j &= \tilde{m}^2\left(c_{e^c}\delta^i_j + d^1_{e^c}\ Y_E^{\alpha\beta i}\left(Y_E^*\right)_{\alpha\beta j}\right)\\
\left(\tilde{m}^2_L\right)^\alpha_\beta &= \tilde{m}^2\left(c_L\delta^\alpha_\beta + d^1_L\ Y_E^{\alpha\gamma k}\left(Y_E^*\right)_{\beta\gamma k} + d^2_L\ Y_D^{\alpha kl}\left(Y_D^*\right)_{\beta kl} + d^3_L\,\mu^\alpha\mu^*_\beta/|\mu|^2\right)\\
B^\alpha &= \tilde{m}^2\left(c_B\,\mu^\alpha/|\mu| + d^1_B\ Y_D^{\alpha kl}(Y_D^*)_{\beta kl}\,\mu^\beta/|\mu| + d^2_B\ Y_E^{\alpha\beta k}(Y_E^*)_{\gamma\beta k}\,\mu^\gamma/|\mu|\right)\\
A_U^{ij} &= A\left(c_{A_U}Y_U^{ij} + d^1_{A_U}Y_U^{kj}\left(Y_D^*\right)_{\alpha kl}Y_D^{\alpha il} + d^2_{A_U}Y_U^{ik}(\lambda''^*)_{klm}(\lambda'')^{jlm}\right.\\
&\quad \left. + d^3_{A_U}Y_U^{ik}\left(Y_U^*\right)_{lk}Y_U^{lj}\right)\\
A_D^{\alpha ij} &= A\left(c_{A_D}Y_D^{\alpha ij} + d^1_{A_D}Y_D^{\alpha kj}\left(Y_U^*\right)_{kl}Y_U^{il} + d^2_{A_D}Y_D^{\beta ij}\left(Y_E^*\right)_{\beta\gamma k}Y_E^{\alpha\gamma k}\right.\\
&\quad + d^3_{A_D}Y_D^{\alpha ik}(\lambda''^*)_{lkm}(\lambda'')^{ljm} + d^4_{A_D}Y_D^{\alpha il}(Y_D^*)_{\gamma kl}(Y_D)^{\gamma kj}\\
&\quad + d^5_{A_D}Y_D^{\alpha kj}(Y_D^*)_{\gamma kl}(Y_D)^{\gamma il} + d^6_{A_D}Y_D^{\alpha kl}(Y_D^*)_{\gamma kl}(Y_D)^{\gamma ij}\\
&\quad \left. + d^7_{A_D}Y_D^{\beta ij}\mu^*_\beta\mu^\alpha/|\mu|^2\right) \qquad\qquad (14)\\
A_E^{\alpha\beta i} &= A\left(c_{A_E}Y_E^{\alpha\beta i} + d^1_{A_E}Y_E^{[\alpha\gamma i}\left(Y_D^*\right)_{\gamma kl}Y_D^{\beta]kl} + d^2_{A_E}Y_E^{\alpha\beta k}\left(Y_E^*\right)_{\gamma\delta k}Y_E^{\gamma\delta i}\right.\\
&\quad \left. + d^3_{A_E}Y_E^{[\alpha\gamma k}\left(Y_E^*\right)_{\gamma\delta k}Y_E^{\beta]\delta i} + d^4_{A_E}Y_E^{[\alpha\gamma i}\mu^*_\gamma\mu^{\beta]}/|\mu|^2\right)\\
A_{\lambda''}^{ijk} &= A\left(c_{A_{\lambda''}}(\lambda'')^{ijk} + d^1_{A_{\lambda''}}(\lambda'')^{ljk}(Y_U^*)_{ml}(Y_U)^{mi}\right.\\
&\quad + d^2_{A_{\lambda''}}(\lambda'')^{i[jl}(Y_D^*)_{\alpha ml}(Y_D)^{\alpha mk]}\\
&\quad \left. + d^3_{A_{\lambda''}}(\lambda'')^{i[jm}(\lambda''^*)_{lnm}(\lambda'')^{lnk]} + d^4_{A_{\lambda''}}(\lambda'')^{imn}(\lambda''^*)_{lmn}(\lambda'')^{ljk}\right)\ ,
\end{aligned}
$$

where the expansion is truncated at the third order in the spurions. The squared brackets stand for anti-symmetrization and we also defined $|\mu|^2 \equiv \sum_{\alpha=1,\ldots,4}|\mu^\alpha|^2$.

In absence of R-parity all the neutral scalar components of $\hat{L}_\alpha$ and $\hat{h}_u$ develop a VEV in order to trigger the electroweak symmetry breaking $SU(2)_L \otimes U(1)_Y \to U(1)_Q$. Given the $SU(4)_{\hat{L}}$ symmetry it is always possible, without loss of generality, to redefine the $\hat{L}_\alpha$ superfield in such a way that only the fourth component acquires a VEV. Then we define operatively the Higgs in such a way that it corresponds to the component which develops a VEV, $\hat{h}_d \equiv \hat{L}_4$, while the leptons do not, $\hat{\ell}_i \equiv \hat{L}_i$.

Despite our notation makes explicit the underlying non-abelian flavour symmetry $SU(3)^4 \otimes SU(4)$, it is also useful to translate it into the more common $SU(3)^5$ language. This connection is provided in Appendix A. Then we can formally split the superpotential in Eq. (8) in an RPC and an RPV term

$$
W_{RPC} = y_U^{ij}\hat{q}_i\hat{u}^c_j\hat{h}_u + y_D^{ij}\hat{h}_d\hat{q}_i\hat{d}^c_j + y_E^{ij}\hat{h}_d\hat{\ell}_i\hat{e}^c_j + \mu\,\hat{h}_u\hat{h}_d\,, \qquad (15)
$$

$$
W_{RPV} = \mu^i\hat{h}_u\hat{\ell}_i + \tfrac{1}{2}\lambda^{ijk}\hat{\ell}_i\hat{\ell}_j\hat{e}^c_k + (\lambda')^{ijk}\hat{\ell}_i\hat{q}_j\hat{d}^c_k + \tfrac{1}{2}(\lambda'')^{ijk}\hat{u}^c_i\hat{d}^c_j\hat{d}^c_k\,, \qquad (16)
$$

and similarly for the soft terms (cf. again Appendix A).

The MFV expansion in Eq. (14) can be easily decomposed in the $SU(3)^5$ language by means of the dictionary given in Eq. (49) of Appendix A

$$
\begin{aligned}
\left(\tilde{m}^2_q\right)^i_j &= \tilde{m}^2 \left(c_q \delta^i_j + d^1_q (y_U y^\dagger_U)^i_j + d^2_q \left[(y_D y^\dagger_D)^i_j + (\lambda')^{lik} \lambda'^*_{ljk}\right]\right) \\
\left(\tilde{m}^2_{u^c}\right)^i_j &= \tilde{m}^2 \left(c_{u^c} \delta^i_j + d^1_{u^c} (y^\dagger_U y_U)^i_j + d^2_{u^c} (\lambda'')^{ikl} (\lambda''^*)_{jkl}\right) \\
\left(\tilde{m}^2_{d^c}\right)^i_j &= \tilde{m}^2 \left(c_{d^c} \delta^i_j + d^1_{d^c} \left[(y^\dagger_D y_D)^i_j + (\lambda')^{lki} \lambda'^*_{lkj}\right] + d^2_{d^c} (\lambda'')^{kil} (\lambda''^*)_{kjl}\right) \\
\left(\tilde{m}^2_{e^c}\right)^i_j &= \tilde{m}^2 \left(c_{e^c} \delta^i_j + d^1_{e^c} \left[2 (y^\dagger_E y_E)^i_j + \lambda^{lki} \lambda^*_{lkj}\right]\right) \\
\left(\tilde{m}^2_\ell\right)^i_j &= \tilde{m}^2 \left(c_L \delta^i_j + d^1_L \left[(y_E y^\dagger_E)^i_j + \lambda^{ilk} \lambda^*_{jlk}\right] + d^2_L (\lambda')^{ilk} \lambda'^*_{jlk} + d^3_L \, \mu^i \mu^*_j / |\mu|^2\right) \\
\left(\tilde{m}^2_d\right)^i &= \tilde{m}^2 \left(d^1_L \lambda^{ilk} (y^*_E)_{lk} + d^2_L (\lambda')^{ilk} (y^*_D)_{lk} + d^3_L \, \mu^i \mu^* / |\mu|^2\right) \\
\tilde{m}^2_{h_d} &= \tilde{m}^2 \left(c_L + d^1_L \ \mathrm{Tr}(y_E y^\dagger_E) + d^2_L \ \mathrm{Tr}(y_D y^\dagger_D) + d^3_L \, \mu \, \mu^* / |\mu|^2\right) \\
b &= \tilde{m}^2 \left(c_B \mu / |\mu| + \ldots\right) \\
b^i &= \tilde{m}^2 \left(c_B \mu^i / |\mu| + \ldots\right) \\
a^{ij}_U &= A \left(c_{A_U} y^{ij}_U + \ldots\right) \\
a^{ij}_D &= A \left(c_{A_D} y^{ij}_D + \ldots\right) \\
a^{ij}_E &= A \left(c_{A_E} y^{ij}_E + \ldots\right) \\
(a_\lambda)^{ijk} &= A \left(c_{A_E} \lambda^{ijk} + \ldots\right) \\
(a_{\lambda'})^{ijk} &= A \left(c_{A_D} (\lambda')^{ijk} + \ldots\right) \\
a^{ijk}_{\lambda''} &= A \left(c_{A_{\lambda''}} (\lambda'')^{ijk} + \ldots\right) \ ,
\end{aligned}
\tag{17}
$$

where for simplicity we have truncated the expansion at the second order in the spurions.

If R-parity is an exact symmetry of the MSSM then neutrinos are massless and there is no lepton flavor violation. Consequently the flavor violation in the lepton sector can be linked to the amount of R-parity violation. For instance the RPV couplings in the expansion of $\tilde{m}^2_\ell$ in Eq. (17) are responsible for flavour violating mass insertions leading to processes like $\ell_i \to \ell_j \gamma$.

Though RPV can be the source of neutrino masses, the MFV expansion is meaningful also in general, being potentially related to not yet measured RPV couplings. However in the next sections we will focus our attention on the case where the neutrino masses are generated by RPV couplings.

3 Neutrino Masses in supersymmetric MFV

In the previous section we formally restored the invariance under the flavour group G_F by promoting all the supersymmetric couplings in Eq. (8) to spurions. Here we want to provide the link between these spurions and the physical observables. Our guideline is to break the flavour group in a minimal way, namely we consider the minimal amount of flavour breaking which is able to reproduce the correct pattern of fermion masses and mixings.

In the limit of massless neutrinos, the connection of the spurions with the flavour structure of the charged fermions is straightforward. From the superpotential

$$W \supset Y_U^{ij}\, \hat{q}_i \hat{u}_j^c \hat{h}_u + Y_D^{4ij}\, \hat{h}_d \hat{q}_i \hat{d}_j^c + Y_E^{4ij}\, \hat{h}_d \hat{\ell}_i \hat{e}_j^c\,, \tag{18}$$

we can identify the relevant flavour violating spurions in terms of known physical observables, up to the parameter $\tan\beta \equiv v_u/v_d$. Indeed it is always possible to choose a basis such that

$$Y_U^{ij} = (V^\dagger \hat{m}_U)^{ij}/v_u\,, \qquad Y_D^{4ij} = \hat{m}_D^{ij}/v_d\,, \qquad Y_E^{4ij} = \hat{m}_E^{ij}/v_d\,, \tag{19}$$

where V is the CKM matrix and $\hat{m}_U$, $\hat{m}_D$, $\hat{m}_E$ are the diagonal charged-fermion masses.

On the other hand the experimental evidence of neutrino masses and mixings makes clear that the flavour group must be further broken. The standard way to introduce neutrino masses in the context of supersymmetric MFV is to extend the field content of the MSSM by introducing three SM-singlet chiral superfields [14, 15] and thus applying the seesaw mechanism [17].

Remarkably the MSSM without R-parity gives the possibility of generating neutrino masses and mixings without the need of additional ingredients. This is the approach we pursue in this work. As we are going to show soon, neutrino masses are fitted by moderate small values of the R-parity violating couplings μ^i/μ, λ and λ', of $\mathcal{O}(10^{-4})$ or even larger. From this point of view the issue of the smallness of neutrino masses could be brought back at the same conceptual level of understanding the flavour structure of the charged fermions, featuring Yukawa couplings also of $\mathcal{O}(10^{-6})$ as in the case of the electron.

The formulae for the neutrino mass matrix in terms of the RPV couplings are collected for completeness in Appendix C. The leading contributions can be schematically written as

$$m_\nu \quad \sim \quad \left(\frac{M_Z}{\tilde{m}}\right)^2 \frac{\mu_i \mu_j}{\tilde{m}}\,, \qquad \frac{3\,\lambda'^2}{8\pi^2}\frac{\hat{m}_D^2}{\tilde{m}}\,, \qquad \frac{\lambda^2}{8\pi^2}\frac{\hat{m}_E^2}{\tilde{m}}\,, \tag{20}$$

where for simplicity we set $M_1 \approx M_2 \approx A \approx \mu \equiv \tilde{m}$ and we neglected the flavour structure of λ' and λ.

For later convenience we split the neutrino mass matrix in a tree level and a one-loop term

$$m_\nu = m_\nu^{(\mathrm{tree})} + m_\nu^{(\mathrm{loop})}\,, \tag{21}$$

whose diagonalization through the PMNS matrix $\hat{U}$ yields

$$m_\nu = \hat{U} \hat{m}_\nu \hat{U}^T\,, \tag{22}$$

where $\hat{m}_\nu$ is the diagonal neutrino mass matrix.

Finally we comment about the baryon number violating coupling λ''. According to our guideline at the beginning of this section, this coupling does not give any contribution to the construction of fermion masses and mixings and thus should be absent as an irreducible source of flavor breaking. However λ'' can still be induced by the other spurions. If the $U(1)$ factors are part of the flavour symmetry that we want to formally restore, then λ'' cannot be generated by the baryon number conserving couplings Y_U, Y_E, Y_D and μ. On the other hand if we consider only the non-abelian symmetry $SU(3)^4 \otimes SU(4)$, the coupling λ'' can be induced in a holomorphic way [15]:

$$\lambda'' \sim Y_U (Y_D)^2 (Y_E)^3 \,, \tag{23}$$

where the proper contractions with the $SU(3)_{\hat{q}}$, $SU(3)_{\hat{e}^c}$ and $SU(4)_{\hat{L}}$ epsilon tensors are understood. Actually, it turns out that the tensor structure forces the invariant to span over RPV couplings and light generation Yukawas, thus providing an automatic suppression of λ''. Remarkably we are able to satisfy the bounds from proton decay without invoking any *ad hoc* conservation or small breaking of the $U(1)_B$ symmetry, but just requiring our minimality condition regarding the identification of the flavor spurions.

3.1 A toy model

When all the RPV spurions are switched on there is an overabundance of free parameters, which cannot all be fixed by the constraints from the neutrino sector. According to our minimality principle we are going to consider scenarios in which only a minimal number of spurions are switched on in order to reproduce neutrino masses and mixings.

Let us illustrate with a toy model of neutrino masses how is it possible to link the spurions with the neutrino observables. Our assumptions are the following:

1. $m_\nu^{(\text{tree})} \gg m_\nu^{(\text{loop})}$

   This readily implies a hierarchical neutrino spectrum. Indeed, in the limit in which only the couplings μ_i are switched on the neutrino mass matrix has rank equal to one, implying only one massive neutrino. Thus we can write

   $$(m_\nu^{(\text{tree})})^{ij} \approx m_3\, \hat{U}^{i3} \hat{U}^{j3} \,, \tag{24}$$

   where m_3 has to satisfy $m_3 \approx \sqrt{\Delta m^2_{\text{atm}}} = 4.9 \cdot 10^{-2}$ eV. Taking $M_1 = M_2 \equiv \tilde{m} \gg M_Z$ in Eq. (59) of Appendix C, we get

   $$\frac{\mu^i}{\mu} = 2.4 \cdot 10^{-5} \left(\frac{\tilde{m}}{1 \text{ TeV}} \right)^{1/2} \left(\frac{\tan\beta}{10} \right) \hat{U}^{i3} \,. \tag{25}$$

2. In order to complete the structure of neutrino mass matrix, we need another source of flavor breaking from the trilinear couplings. Among these couplings we turn on,

as an example, only λ'^{i33}. Our assumption is that these couplings are completely responsible for the second neutrino. This implies

$$(m_\nu^{(\text{loop})})^{ij} \approx m_2\, \hat{U}^{i2}\hat{U}^{j2}\,, \tag{26}$$

where $m_2 \gtrsim \sqrt{\Delta m^2_{\text{sol}}} = 8.7 \cdot 10^{-3}$ eV. At the leading order in the MFV expansion and taking $c_{A_D} = c_{d^c} = c_q = 1$ (cf. Eq. (63) in Appendix C), we get

$$(\lambda')^{i33} = 3.3 \cdot 10^{-5} \left(\frac{m_2}{8.7 \cdot 10^{-3}\ \text{eV}}\right)^{1/2} \left(\frac{\tilde{m}}{1\ \text{TeV}}\right)^{1/2} \left(\frac{\tan\beta}{10}\right)^{1/2} \hat{U}^{i2}\,. \tag{27}$$

Once the relevant spurions are fixed in terms of the neutrino masses and mixings one can use the MFV expansion in order to make predictions for LFV processes.

In order to properly determine LFV processes, one as to consider several kind of contributions (see [18] for an example of computation). Since we are interested in an order of magnitude estimate of the processes induced in our MLFV setup, we will just focus on the effects induced by the non-diagonal entries in the sfermion mass matrices due to the spurions. In this case the normalized branching ratios for the processes $\ell_i \to \ell_j\gamma$ are given by [19]:

$$\frac{BR(\ell_i \to \ell_j\,\gamma)}{BR(\ell_i \to \ell_j\nu_i\bar{\nu}_j)} \approx \frac{\alpha^3}{G_F^2}\frac{\delta_{ij}^2}{\tilde{m}^4}\tan^2\beta\,, \tag{28}$$

where the flavour violating mass insertion δ_{ij} can be expressed as combinations of neutrino masses and elements of the PMNS matrix, according to the MFV expansion. For instance in our toy model where only the couplings μ^i and λ'^{i33} are switched on, δ^{LL}_{ij} reads (cf. Eq. (17))

$$\delta^{LL}_{ij} = \frac{\Delta^{LL}_{ij}}{\tilde{m}^2_\ell} \approx \frac{1}{c_L}\left(d_L^2(\lambda')^{i33}\lambda'^*_{j33} + d_L^3\,\frac{\mu^i\mu^*_j}{|\mu|^2}\right)\,, \tag{29}$$

where Δ^{LL}_{ij} is the flavour violating part of $\tilde{m}^2_\ell$. As it is evident from Eq. (29), the mass insertions scale like the square of the RPV parameters. Given the following estimation of the branching ratios in Eq. (28)

$$\frac{BR(\ell_i \to \ell_j\,\gamma)}{BR(\ell_i \to \ell_j\nu_i\bar{\nu}_j)} \approx 10^{-27}\left(\frac{\tilde{m}}{1\text{TeV}}\right)^{-4}\left(\frac{\tan\beta}{10}\right)^2\left(\frac{\lambda'}{10^{-5}}\right)^4\,, \tag{30}$$

one concludes that it is not possible to accomplish observable rates, in view of the current experimental bounds showed in Table 1.

Let us mention that rates of $\mu \to e\,\gamma$ closer to the experimental sensitivity can be obtained when neutrino masses are fitted by trilinears featuring first families indices, like for instance λ'^{i11}. In such a case the suppression due to the down-quark mass in the expression of the neutrino mass matrix (cf. Eq. (63) in Appendix C) allows for larger values of λ'^{i11} even of $\mathcal{O}(10^{-2})$. However such a large coupling may be in conflict with

LFV process	current bound
$BR(\mu \to e\,\gamma)$	2.4×10^{-12} [20]
$BR(\tau \to e\,\gamma)$	1.1 $\times 10^{-7}$ [21]
$BR(\tau \to \mu\,\gamma)$	4.5 $\times 10^{-8}$ [22]
$BR(\mu \to 3\,e)$	1.1 $\times 10^{-11}$ [23]
$BR(\tau \to 3\,e)$	3.6 $\times 10^{-8}$ [23]
$BR(\tau \to 3\,\mu)$	3.2 $\times 10^{-8}$ [23]
$CR(\mu \to e\,, Ti)$	4.3×10^{-12} [23]
$CR(\mu \to e\,, Au)$	7.0×10^{-12} [23]

Table 1: Summary of the current experimental bounds on LFV processes. For later convenience we reported also the current bounds on $\ell_i \to \ell_j \ell_k \ell_k$ decays and $\mu \to e$ conversions in nuclei.

other flavour violating observables [12]. A complete analysis of such scenarios and a more realistic model for neutrino masses is postponed to future works.

In the next section we are going to consider another symmetry pattern in which the breaking of the lepton number is separated from that of lepton flavour. In this setup the mass insertions are enhanced by a factor $1/\varepsilon_L^2$, where ε_L is related to the amount of breaking of the total lepton number, thus allowing to lift the rates towards the experimental sensitivity.

4 A predictive scenario: $SU(3)^5 \otimes U(1)_L \otimes U(1)_B$

In the previous section we have seen that the contribution to LFV processes are generically well below the present experimental bounds. This is due to fact that the spurions responsible for neutrino masses break simultaneously both the total lepton number and the non-abelian part of the flavor group. As it has been shown in [4, 9], in order to have measurable rates for the flavor changing radiative charged lepton decays, one has to separate the source of breaking of lepton number from that of LFV.

This leads us to consider a different scenario based on another subgroup of the original kinetic symmetry $G_{\rm kin}^{\rm MSSM}$ (cf. Eq. (5)). We assume that the symmetry that we want to formally restore is given by $SU(3)^5 \otimes U(1)_L \otimes U(1)_B$, where L and B are the total lepton and baryon number. The $U(1)_B$ factor is needed in order to properly suppress dangerous contributions to the proton decay rate (for the relevant bounds see for instance Ref. [24]).

In this setup the R-parity violating couplings μ_i, λ, λ' and λ'' can be split in two parts, one responsible for the breaking of lepton and baryon number and the other for the breaking of the flavor group

$$\mu^i = \varepsilon_L \tilde{\mu}^i\,, \qquad \lambda = \varepsilon_L \tilde{\lambda}\,, \qquad \lambda' = \varepsilon_L \tilde{\lambda}'\,, \qquad \lambda'' = \varepsilon_B \tilde{\lambda}''\,. \tag{31}$$

The quantum numbers of the flavor spurions under $SU(3)^5$ are given by

$$\begin{aligned} y_U &\sim (\bar{3},\bar{3},1,1,1) \\ y_D &\sim (\bar{3},1,\bar{3},1,1) \\ y_E &\sim (1,1,1,\bar{3},\bar{3}) \\ \tilde{\mu} &\sim (1,1,1,1,\bar{3}) \\ \tilde{\lambda} &\sim (1,1,1,\bar{3},3) \\ \tilde{\lambda}' &\sim (\bar{3},1,\bar{3},1,\bar{3}) \\ \tilde{\lambda}'' &\sim (1,\bar{3},3,1,1)\,, \end{aligned} \tag{32}$$

while ε_L and ε_B have charge -1 and $+1$ respectively under $U(1)_L$ and $U(1)_B$. The corresponding MFV expansion is reported in Eq. (54) of Appendix B.2.

In this case the rates of the LFV processes are dominated by the lepton and baryon number preserving (but flavor changing) slepton mass insertions, which depend only on the parameters $\tilde{\mu}$, $\tilde{\lambda}$, $\tilde{\lambda}'$. Other RPV vertex contributions depend on quantities which violate total lepton number and hence are suppressed by the ε_L factor.

As we are going to show, peculiar correlations among physical observables will emerge due the MFV expansion. For definiteness we consider an example based on the toy model of neutrino masses already introduced in the previous section, where only the spurions μ^i and λ'^{i33} are switched on.

In such a case the relevant off-diagonal terms $i \neq j$ induced by this two spurions in $(\tilde{m}^2_\ell)^i_j$ and a^{ij}_E are given by (cf. Eq. (54) in Appendix B)

$$\left(\tilde{m}^2_\ell\right)^i_j = \tilde{m}^2 \left(d^2_\ell (\tilde{\lambda}')^{i33} \tilde{\lambda}'^*_{j33} + d^3_\ell \frac{\tilde{\mu}^i \tilde{\mu}^*_j}{|\mu|^2} \right)\,, \tag{33}$$

$$a^{ij}_E = A\, y^{jj}_E \left(d^4_{a_E} \frac{\tilde{\mu}^*_j \tilde{\mu}^i}{|\mu|^2} + d^5_{a_E} \tilde{\lambda}'^*_{j33} (\tilde{\lambda}')^{i33} \right)\,. \tag{34}$$

In our setup it turns out that the LL mass insertions, and thus $(\tilde{m}^2_\ell)^i_j$, give the dominant contribution to the LFV processes[1]. Focusing on δ^{LL}, using (24) and (26), we get:

$$\begin{aligned}
\left(\delta^{LL}\right)^i_j &= \frac{1}{c_\ell}\left[d_\ell^3 \frac{\tilde{\mu}^i \tilde{\mu}^*_j}{|\mu|^2} + d_\ell^2 (\tilde{\lambda}')^{i33} \tilde{\lambda}'^*_{j33}\right] \\
&= \frac{1}{\varepsilon_L^2 c_\ell}\left[d_\ell^3 \left(\frac{\tan\beta}{M_Z}\right)^2 \left(\frac{M_1 M_2}{M_1 c_W^2 + M_2 s_W^2} - \frac{M_Z^2}{\mu}\sin 2\beta\right) m_3\, \hat{U}^{i3}(\hat{U}^{j3})^* \right. \\
&\qquad \left. + d_\ell^2 \frac{8\pi^2 \tilde{m}^2}{3\mu \tan\beta\, m_b^2}\, m_2\, \hat{U}^{i2}(\hat{U}^{j2})^*\right]. \qquad (35)
\end{aligned}$$

Notice that the factor $1/\varepsilon_L^2$ in the mass insertions implies an enhancement of $1/\varepsilon_L^4$ in the rates. Indeed it is possible to estimate the branching ratios in the following way

$$\frac{BR(\ell_i \to \ell_j\, \gamma)}{BR(\ell_i \to \ell_j \nu_i \overline{\nu}_j)} \approx 10^{-27} \left(\frac{1}{\varepsilon_L}\right)^4 \left(\frac{\tilde{m}}{1\ \text{TeV}}\right)^{-4} \left(\frac{\tan\beta}{10}\right)^2 \left(\frac{\lambda'}{10^{-5}}\right)^4, \qquad (36)$$

for values of $\varepsilon_L \sim 10^{-(3\div 4)}$ the rates of the three relevant processes can get close to the experimental sensitivities, depending on the values of SUSY parameters.

Notice that the coupling ε_L cannot be arbitrarily small. Indeed, by imposing the relations in Eqs. (25)–(27) and by requiring that the flavor violating parameters $\tilde{\mu}$ and $\tilde{\lambda}'$ are at most of order one, we can estimate the lower bound $\varepsilon_L \gtrsim 10^{-5}$.

Given the potential detectability of these processes it is now interesting to compute the ratio among the branching ratios of the LFV channels. The correlation of these quantities with the neutrino observables allows for peculiar predictions, hence giving a way to distinguish our realization of MFV from other setups. In our model of neutrino masses the ratio between two branching ratios is given by the ratio of the mass insertions squared and it is parametrized by

$$\frac{BR(\ell_j \to \ell_i \gamma)}{BR(\ell_k \to \ell_m \gamma)} = \frac{|\hat{U}^{i2}(\hat{U}^{j2})^* + c\, \hat{U}^{i3}(\hat{U}^{j3})^*|^2}{|\hat{U}^{m2}(\hat{U}^{k2})^* + c\, \hat{U}^{m3}(\hat{U}^{k3})^*|^2}, \qquad (37)$$

where the constant c can be estimated as

$$c \approx 1.4 \times 10^{-1} \left(\frac{d_\ell^3}{d_\ell^2}\right) \left(\frac{\tan\beta}{10}\right)^3 \left(\frac{\mu}{1\ \text{TeV}}\right) \left(\frac{\tilde{m}}{1\ \text{TeV}}\right)^{-2} \left(\frac{M_G}{300\ \text{GeV}}\right), \qquad (38)$$

where $M_1 = M_2 \equiv M_G$ and we have imposed $m_2 = \sqrt{\Delta m^2_{\text{sol}}}$ and $m_3 = \sqrt{\Delta m^2_{\text{atm}}}$. From Eq. (38) it is evident that, depending on the SUSY parameters, the mass insertions are dominated either by the trilinear ($c \ll 1$) or the bilinear ($c \gg 1$) couplings. It is possible then to identify two asymptotic regimes in which the ratio between the LFV branching ratios have a simple analitical expression:

[1] The term a_E is responsible for the LR mass insertions. However, according to the analysis of Refs. [19], δ^{LR} is negligible provided that $\delta^{LR}_{ij} \ll (m_{\ell_i}/\tilde{m}) \tan\beta\ \delta^{LL}_{ij}$. In our case, assuming all the coefficients of the MFV expansion to be of order one, this condition translates into $|(\tilde{m}/A)\tan\beta| \gg 1$.

- $|c| \ll 1$

  In this case the mass insertions are dominated by the contribution from the trilinear couplings $(\tilde{\lambda}')^{i33}$. Since the ratios $BR(\mu \to e\gamma)/BR(\tau \to \mu\gamma)$ and $BR(\mu \to e\gamma)/BR(\tau \to e\gamma)$ show only a slight dependence from the Maiorana phase δ, we take $\delta = 0, \pi$ and obtain

$$\frac{BR(\mu \to e\gamma)}{BR(\tau \to \mu\gamma)} \approx \frac{|\hat{U}^{12}|^2}{|\hat{U}^{32}|^2} = \frac{s_{12}^2 c_{13}^2}{(\mp c_{23}s_{12}s_{13} - c_{12}s_{23})^2} \approx 0.53 \div 1.75\,, \tag{39}$$

$$\frac{BR(\mu \to e\gamma)}{BR(\tau \to e\gamma)} \approx \frac{|\hat{U}^{22}|^2}{|\hat{U}^{32}|^2} = \frac{(c_{12}c_{23} \mp s_{12}s_{13}s_{23})^2}{(\mp c_{23}s_{12}s_{13} - c_{12}s_{23})^2} \approx 0.37 \div 2.4\,, \tag{40}$$

  where the extrema of the range are obtained by scanning over the 2-σ values of the mixing angles (cf. Table 2). In this case, the three branching ratios are of the same order of magnitude. Notice that the LFV effects depend on the PMNS matrix $\hat{U}^{i2}$, differently with respect to other MLFV setups (cf. for instance Table 3).

- $|c| \gg 1$

  In this case the mass insertions are dominated by the bilinear couplings μ^i. Then we can derive the following functional behaviors for the two relevant ratios

$$\frac{BR(\mu \to e\gamma)}{BR(\tau \to \mu\gamma)} \approx \frac{|\hat{U}^{13}|^2}{|\hat{U}^{33}|^2} \approx \frac{s_{13}^2}{c_{13}^2 c_{23}^2} \approx 0.007 \div 0.07\,, \tag{41}$$

$$\frac{BR(\mu \to e\gamma)}{BR(\tau \to e\gamma)} \approx \frac{|\hat{U}^{23}|^2}{|\hat{U}^{33}|^2} \approx \frac{s_{23}^2}{c_{23}^2} \approx 0.7 \div 1.6\,. \tag{42}$$

  Compared to the previous case we observe an enhancement of $BR(\tau \to \mu\,\gamma)$ compared to the other branching ratios. This results coincides with the one found in Ref. [11] in the case of inverted hierarchy of neutrino masses.

Furthermore, in order to study the general case we vary the parameter c in the range $[-100, 100]$ and the parameters of the neutrino sector according to Table 2 as before. The results are plotted in Fig. 1. In Fig. 2 we report the correlation between the two ratios $BR(\mu \to e\,\gamma)/BR(\tau \to \mu\,\gamma)$ and $BR(\mu \to e\,\gamma)/BR(\tau \to e\,\gamma)$.

From Fig. 1 it is possible to see that, away from the two asymptotic regimes, there are regions of strong enhancement or suppression of the ratios. Indeed we can estimate the values of the c parameter for which $BR(\ell_i \to \ell_j\,\gamma) \to 0$

$$BR(\mu \to e\,\gamma) \to 0 \qquad \longrightarrow \qquad c \approx \mp\frac{c_{12}s_{23}s_{12}}{s_{23}s_{13}} = \mp(2.23 \div 9.43) \qquad (\delta = 0, \pi)\,, \tag{43}$$

$$BR(\tau \to e\,\gamma) \to 0 \qquad \longrightarrow \qquad c \approx \pm\frac{c_{12}s_{23}s_{12}}{s_{23}s_{13}} = \pm(2.23 \div 9.43) \qquad (\delta = 0, \pi)\,, \tag{44}$$

Observable	Best fit	2-σ
$\Delta m^2_{\rm atm}$	2.50×10^{-3} eV^2	$(2.25 - 2.68) \times 10^{-3}$ eV^2
$\Delta m^2_{\rm sol}$	7.59×10^{-5} eV^2	$(7.24 - 7.99) \times 10^{-5}$ eV^2
$\sin^2\theta_{12}$	0.312	$0.28 - 0.35$
$\sin^2\theta_{23}$	0.52	$0.41 - 0.61$
$\sin^2\theta_{13}$	0.013	$0.004 - 0.028$

Table 2: Experimental values of the neutrino sector observables as reported in Ref. [25]. For the PMNS matrix we have considered the PDG parametrization [23]. The Dirac phase δ varies in the range $[0, 2\pi]$.

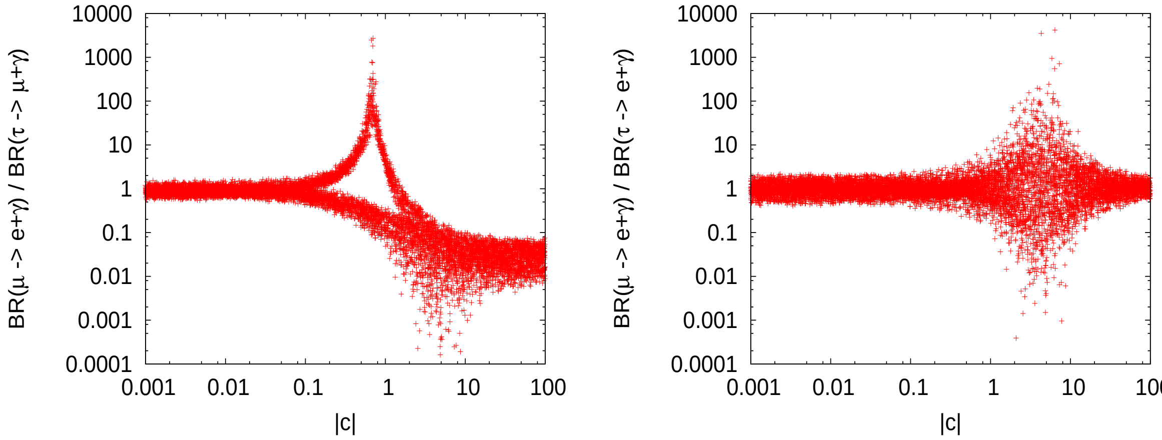

Figure 1: Ratios between branching ratios of as function of $|c|$.

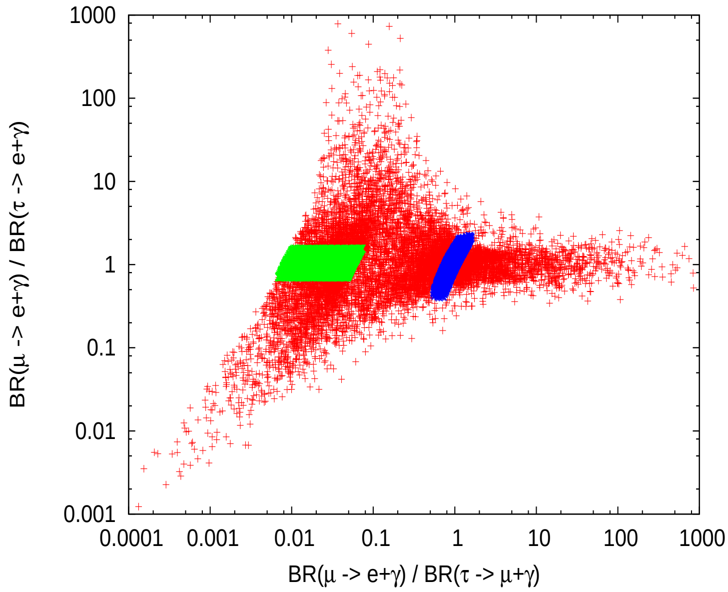

Figure 2: $BR(\mu \to e\,\gamma)/BR(\tau \to \mu\,\gamma)$ versus $BR(\mu \to e\,\gamma)/BR(\tau \to e\,\gamma)$. The parameter c is varied in the range $[-100, 100]$. The blue region is characterized by $|c| = 0$ while the green one by $|c| \gg 1$.

$$BR(\tau \to \mu\,\gamma) \to 0 \qquad \longrightarrow \qquad c \approx \frac{c_{12}^2}{c_{13}^2} = (0.66 \div 0.72)\,. \tag{45}$$

Finally, for the purpose of comparison, we show in Table 3 the LFV parameters predicted by various MLFV models.

Model	Flavor violating parameter
Minimal field content [4]	$\hat{U}\hat{m}_\nu^2\hat{U}^\dagger$
Extended field content + CP limit [4]	$\hat{U}\hat{m}_\nu\hat{U}^\dagger$
Extended field content + leptogenesis [6, 8]	$\hat{U}\hat{m}_\nu^{1/2}\,H^2\,\hat{m}_\nu^{1/2}\hat{U}^\dagger$
$SU(3)_\ell \otimes SU(3)_N \to SU(3)_{\ell+N}$ [11]	$\hat{U}\frac{1}{\hat{m}_\nu^2}\hat{U}^\dagger$
MSSM wihtout R-parity (toy model)	$\hat{U}^{i2}\hat{U}^{*j2} + c\,\hat{U}^{i3}\hat{U}^{*j3}$

Table 3: Comparative summary of MLFV models.

We conclude commenting on other LFV processes like $\mu \to e$ conversions and $\ell_i \to \ell_j\ell_k\ell_k$ decays, not considered until now. In our case these processes are determined by γ-penguin type diagrams [26] and turn out to have the same flavor structure of $\ell_i \to \ell_j\gamma$. This implies in particular that the decays in three leptons have similar patterns of enhancements/suppressions of those discussed above. Notice however that the radiative decays are the processes most severely constrained by the experiments.

5 Conclusions

In this work we presented a general supersymmetric version of MFV including also the RPV terms as the irreducible sources of the flavour symmetry breaking. If the RPV couplings are responsible for neutrino masses, the framework can be viewed as an extension of MFV to the lepton sector.

An important aspect stressed throughout the paper is that the global symmetry of the kinetic term of the MSSM lagrangian is enhanced with respect that of the SM. Indeed the superfields $\hat{\ell}_i$ and $\hat{h}_d$ can be rearranged in a 4-dimensional flavour multiplet $\hat{L}$, whose kinetic term is invariant under $U(4)_{\hat{L}}$ unitary transformations. This gives us the possibility to consider as the most general flavour symmetry the non-abelian group $SU(3)^4 \otimes SU(4)_{\hat{L}}$. In such a case the breaking of the total lepton number and that of lepton flavour number are linked together, thus generically implying small effects in LFV physics.

On the other hand the separation between the breaking of lepton number and lepton flavour number leads to an interesting phenomenology. This is the motivation to consider our second scenario based on the $SU(3)^5 \otimes U(1)_L \otimes U(1)_B$ flavour symmetry. This last option yields peculiar correlations among the branching ratios of the $\ell_i \to \ell_j \gamma$ processes.

Several interesting possibilities could be taken into account for future investigations both from a theoretical and a phenomenological point of view. For instance, the neutrino mass model employed here should be considered just as a toy model which allows to easily connect the RPV spurions with the observables in the neutrino sector, and a more realistic model should be taken into account.

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A Notation

In this Appendix we define both the $SU(3)^4 \otimes SU(4)$ and $SU(3)^5$ notations and provide the translation between the two languages.

- $SU(3)^4 \otimes SU(4)$ notation

$$W = Y_U^{ij} \hat{q}_i \hat{u}_j^c \hat{h}_u + Y_D^{\alpha ij} \hat{L}_\alpha \hat{q}_i \hat{d}_j^c + \tfrac{1}{2} Y_E^{\alpha\beta i} \hat{L}_\alpha \hat{L}_\beta \hat{e}_i^c + \mu^\alpha \hat{h}_u \hat{L}_\alpha + \tfrac{1}{2} (\lambda'')^{ijk} \hat{u}_i^c \hat{d}_j^c \hat{d}_k^c$$

$$\begin{aligned} -\mathcal{L}_{\text{soft}} = & \text{ gaugino masses} \\ + & \sum_F \tilde{F}^\dagger \tilde{m}_F^2 \tilde{F} + \tilde{m}_{h_u}^2 h_u^* h_u + \left(B^\alpha h_u \tilde{L}_\alpha + \text{h.c.} \right) \\ + & A_U^{ij} \tilde{q}_i \tilde{u}_j^c h_u + A_D^{\alpha ij} \tilde{L}_\alpha \tilde{q}_i \tilde{d}_j^c + \tfrac{1}{2} A_E^{\alpha\beta j} \tilde{L}_\alpha \tilde{L}_\beta \tilde{e}_j^c + \tfrac{1}{2} (A_{\lambda''})^{ijk} \tilde{u}_i^c \tilde{d}_j^c \tilde{d}_k^c + \text{h.c.} \end{aligned} \tag{46}$$

  with $F = \{q, u^c, d^c, e^c, L\}$.

- $SU(3)^5$ notation

$$\begin{aligned} W_{\text{RPC}} = & \; y_U^{ij} \hat{q}_i \hat{u}_j^c \hat{h}_u + y_D^{ij} \hat{h}_d \hat{q}_i \hat{d}_j^c + y_E^{ij} \hat{h}_d \hat{\ell}_i \hat{e}_j^c + \mu \, \hat{h}_u \hat{h}_d \\ W_{\text{RPV}} = & \; \mu^i \hat{h}_u \hat{\ell}_i + \tfrac{1}{2} \lambda^{ijk} \hat{\ell}_i \hat{\ell}_j \hat{e}_k^c + (\lambda')^{ijk} \hat{\ell}_i \hat{q}_j \hat{d}_k^c + \tfrac{1}{2} (\lambda'')^{ijk} \hat{u}_i^c \hat{d}_j^c \hat{d}_k^c \\ -\mathcal{L}_{\text{soft}}^{\text{RPC}} = & \text{ gaugino masses} \\ + & \sum_f \tilde{f}^\dagger \tilde{m}_f^2 \tilde{f} + \tilde{m}_{h_u}^2 h_u^* h_u + \tilde{m}_{h_d}^2 h_d^* h_d + (b \, h_u h_d + \text{h.c.}) \\ + & \; a_U^{ij} \tilde{q}_i \tilde{u}_j^c h_u + a_D^{ij} h_d \tilde{q}_i \tilde{d}_j^c + a_E^{ij} h_d \tilde{\ell}_i \tilde{e}_j^c + \text{h.c.} \\ -\mathcal{L}_{\text{soft}}^{\text{RPV}} = & \; (\tilde{m}_d^2)^i h_d^* \tilde{\ell}_i + b^i h_u \tilde{\ell}_i + \text{h.c.} \\ + & \tfrac{1}{2} (a_\lambda)^{ijk} \tilde{\ell}_i \tilde{\ell}_j \tilde{e}_k^c + (a_{\lambda'})^{ijk} \tilde{\ell}_i \tilde{q}_j \tilde{d}_k^c + \tfrac{1}{2} (a_{\lambda''})^{ijk} \tilde{u}_i^c \tilde{d}_j^c \tilde{d}_k^c + \text{h.c.} \end{aligned} \tag{47}$$

  with $f = \{q, u^c, d^c, e^c, \ell\}$.

Let us define

$$\hat{L}_i \equiv \hat{\ell}_i \,, \qquad \hat{L}_4 \equiv \hat{h}_d \,, \qquad \tilde{L}_4 \equiv h_d \,, \tag{48}$$

then, by comparing Eq. (47) with Eq. (46), the following identifications follow

$$
\begin{aligned}
y_U^{ij} &= Y_U^{ij} \\
y_D^{ij} &= Y_D^{4ij} \\
y_E^{ij} &= Y_E^{4ij} \\
\mu &= \mu^4 \\
\lambda^{ijk} &= Y_E^{ijk} \\
(\lambda')^{ijk} &= Y_D^{ijk} \\
(\tilde{m}_\ell)^i_j &= (\tilde{m}_L)^i_j \\
\tilde{m}_{h_d} &= (\tilde{m}_L)^4_4 \\
(\tilde{m}_d)^i &= (\tilde{m}_L)^i_4 \\
b &= B^4 \\
b^i &= B^i \\
a_U^{ij} &= A_U^{ij} \\
a_D^{ij} &= A_D^{4ij} \\
a_E^{ij} &= A_E^{4ij} \\
(a_\lambda)^{ijk} &= A_E^{ijk} \\
(a_{\lambda'})^{ijk} &= A_D^{ijk}\,.
\end{aligned}
\tag{49}
$$

B Group theory

B.1 $SU(3)^4 \otimes SU(4)$

Spurions:

$$
\begin{aligned}
\mu &\sim (1,1,1,1,\bar{4}) \\
Y_U &\sim (\bar{3},\bar{3},1,1,1) \\
Y_D &\sim (\bar{3},1,\bar{3},1,\bar{4}) \\
Y_E &\sim (1,1,1,\bar{3},6) \\
\lambda'' &\sim (1,\bar{3},3,1,1)
\end{aligned}
\tag{50}
$$

Soft terms:

$$
\begin{aligned}
\tilde{m}_q^2 &\sim (8,1,1,1,1) \oplus (1,1,1,1,1) \\
\tilde{m}_{u^c}^2 &\sim (1,8,1,1,1) \oplus (1,1,1,1,1) \\
\tilde{m}_{d^c}^2 &\sim (1,1,8,1,1) \oplus (1,1,1,1,1) \\
\tilde{m}_{e^c}^2 &\sim (1,1,1,8,1) \oplus (1,1,1,1,1) \\
\tilde{m}_L^2 &\sim (1,1,1,1,15) \oplus (1,1,1,1,1) \\
B &\sim (1,1,1,1,\bar{4}) \\
A_U &\sim (\bar{3},\bar{3},1,1,1) \\
A_D &\sim (\bar{3},1,\bar{3},1,\bar{4}) \\
A_E &\sim (1,1,1,\bar{3},6)
\end{aligned}
\tag{51}
$$

The expansion of the soft terms in both the $SU(3)^4 \otimes SU(4)$ and the $SU(3)^5$ languages is provided respectively in Eq. (14) and Eq. (17) of Sect. 2.

B.2 $SU(3)^5 \otimes U(1)_L \otimes U(1)_B$

Spurions:

$$\begin{array}{rcl} y_U & \sim & (\bar{3},\bar{3},1,1,1)_{(0,0)} \\ y_D & \sim & (\bar{3},1,\bar{3},1,1)_{(0,0)} \\ y_E & \sim & (1,1,1,\bar{3},\bar{3})_{(0,0)} \\ \tilde{\mu}^i & \sim & (1,1,1,1,\bar{3})_{(0,0)} \\ \tilde{\lambda} & \sim & (1,1,1,\bar{3},3)_{(0,0)} \\ \tilde{\lambda}' & \sim & (\bar{3},1,\bar{3},1,\bar{3})_{(0,0)} \\ \tilde{\lambda}'' & \sim & (1,\bar{3},3,1,1)_{(0,0)} \\ \varepsilon_L & \sim & (1,1,1,1,1)_{(-1,0)} \\ \varepsilon_B & \sim & (1,1,1,1,1)_{(0,+1)} \end{array} \tag{52}$$

where the subscripts label the abelian quantum numbers.

Soft terms:

$$\begin{array}{rcl} \tilde{m}_q^2 & \sim & (8,1,1,1,1)_{(0,0)} \oplus (1,1,1,1,1)_{(0,0)} \\ \tilde{m}_{u^c}^2 & \sim & (1,8,1,1,1)_{(0,0)} \oplus (1,1,1,1,1)_{(0,0)} \\ \tilde{m}_{d^c}^2 & \sim & (1,1,8,1,1)_{(0,0)} \oplus (1,1,1,1,1)_{(0,0)} \\ \tilde{m}_{e^c}^2 & \sim & (1,1,1,8,1)_{(0,0)} \oplus (1,1,1,1,1)_{(0,0)} \\ \tilde{m}_l^2 & \sim & (1,1,1,1,8)_{(0,0)} \oplus (1,1,1,1,1)_{(0,0)} \\ (\tilde{m}_d^2)^i & \sim & (1,1,1,1,\bar{3})_{(+1,0)} \\ b^i & \sim & (1,1,1,1,\bar{3})_{(+1,0)} \\ a_U & \sim & (\bar{3},\bar{3},1,1,1)_{(0,0)} \\ a_D & \sim & (\bar{3},1,\bar{3},1,1)_{(0,0)} \\ a_E & \sim & (1,1,1,\bar{3},\bar{3})_{(0,0)} \\ a_\lambda & \sim & (1,1,1,\bar{3},3)_{(+1,0)} \\ a_{\lambda'} & \sim & (\bar{3},1,\bar{3},1,\bar{3})_{(+1,0)} \\ a_{\lambda''} & \sim & (1,\bar{3},3,1,1)_{(0,-1)} \end{array} \tag{53}$$

Then the MFV expansion reads

$$
\begin{aligned}
\left(\tilde{m}_q^2\right)^i_{\ j} &= \tilde{m}^2\left(c_q\delta^i_j + d_q^1(y_U y_U^\dagger)^i_j + d_q^{(21)}(y_D y_D^\dagger)^i_j + d_q^{(22)}(\tilde{\lambda}')^{lik}\tilde{\lambda}'^*_{ljk}\right)\\
\left(\tilde{m}_{u^c}^2\right)^i_{\ j} &= \tilde{m}^2\left(c_{u^c}\delta^i_j + d_{u^c}^1(y_U^\dagger y_U)^i_j + d_{u^c}^2(\tilde{\lambda}'')^{ikl}(\tilde{\lambda}''^*)_{jkl}\right)\\
\left(\tilde{m}_{d^c}^2\right)^i_{\ j} &= \tilde{m}^2\left(c_{d^c}\delta^i_j + d_{d^c}^{(11)}(y_D^\dagger y_D)^i_j + d_{d^c}^{(12)}(\tilde{\lambda}')^{lki}\tilde{\lambda}'^*_{lkj} + d_{d^c}^2(\tilde{\lambda}'')^{kil}(\tilde{\lambda}''^*)_{kjl}\right)\\
\left(\tilde{m}_{e^c}^2\right)^i_{\ j} &= \tilde{m}^2\left(c_{e^c}\delta^i_j + d_{e^c}^{(11)}(y_E^\dagger y_E)^i_j + d_{e^c}^{(12)}\tilde{\lambda}^{lki}\tilde{\lambda}^*_{lkj}\right)\\
\left(\tilde{m}_{\ell}^2\right)^i_{\ j} &= \tilde{m}^2\left(c_\ell\delta^i_j + d_\ell^{(11)}(y_E y_E^\dagger)^i_j + d_\ell^{(12)}\tilde{\lambda}^{ilk}\tilde{\lambda}^*_{jlk} + d_\ell^2(\tilde{\lambda}')^{ilk}\tilde{\lambda}'^*_{jlk} + d_\ell^3\,\tilde{\mu}^i\tilde{\mu}^*_j/|\mu|\right)\\
\left(\tilde{m}_d^2\right)^i &= \tilde{m}^2\,\varepsilon_L\left(d_d^1\,\tilde{\mu}^i/|\mu| + d_d^2\,(\tilde{\lambda}')^{ilk}(y_D^*)_{lk} + d_d^3\,(\tilde{\lambda}')^{ilk}\tilde{\lambda}'^*_{plk}\,\tilde{\mu}^p/|\mu|\right.\\
&\qquad \left. + d_d^4\,(\tilde{\lambda})^{ilk}(y_E^*)_{lk} + d_d^5\,(\tilde{\lambda})^{ilk}\tilde{\lambda}^*_{plk}\,\tilde{\mu}^p/|\mu| + d_d^6\,(y_E y_E^\dagger)^i_p\,\tilde{\mu}^p/|\mu|\right)\\
b^i &= \tilde{m}^2\,\varepsilon_L\left(c_b\,\tilde{\mu}^i/|\mu| + d_b^{(11)}(\tilde{\lambda}')^{ilk}(y_D^*)_{lk} + d_b^{(12)}(\tilde{\lambda}')^{ilk}\tilde{\lambda}'^*_{plk}\,\tilde{\mu}^p/|\mu|\right.\\
&\qquad \left. + d_b^{(21)}(\tilde{\lambda})^{ilk}(y_E^*)_{lk} + d_b^{(22)}(\tilde{\lambda})^{ilk}\tilde{\lambda}^*_{plk}\,\tilde{\mu}^p/|\mu| + d_b^{(23)}(y_E y_E^\dagger)^i_p\,\tilde{\mu}^p/|\mu|\right)\\
a_U^{ij} &= A\left(c_{a_U}y_U^{ij} + \ldots\right)\\
a_D^{ij} &= A\left(c_{a_D}y_D^{ij} + \ldots\right)\\
a_E^{ij} &= A\left(c_{a_E}y_E^{ij} + d_{a_E}^1(y_E y_E^\dagger y_E)^{ij} + d_{a_E}^2\tilde{\lambda}^{ikj}\tilde{\lambda}'^*_{klm}y_D^{lm} + d_{a_E}^3 y_E^{kj}\tilde{\mu}^*_k\tilde{\mu}^i/|\mu|\right.\\
&\qquad + d_{a_E}^4 y_E^{kj}\tilde{\lambda}'^*_{klm}(\tilde{\lambda}')^{ilm} + d_{a_E}^5 y_E^{im}\tilde{\lambda}^*_{klm}\tilde{\lambda}^{klj} + d_{a_E}^6 y_E^{kj}\tilde{\lambda}^*_{klm}\tilde{\lambda}^{ilm}\\
&\qquad + d_{a_E}^7 y_E^{km}\tilde{\lambda}^*_{klm}\tilde{\lambda}^{lij} + d_{a_E}^8\epsilon_{klm}\tilde{\lambda}^{ikj}\tilde{\mu}^l\tilde{\mu}^m + d_{a_E}^9\epsilon^{ikl}\epsilon^{jmn}(y_E^*)_{kl}(y_E^*)_{mn}\\
&\qquad + d_{a_E}^{10}\epsilon^{ikl}\epsilon^{jmn}\tilde{\mu}^p(y_E^*)_{pm}\tilde{\lambda}^*_{kln} + d_{a_E}^{11}\epsilon^{ikl}\epsilon^{jmn}\tilde{\mu}^p(y_E^*)_{km}\tilde{\lambda}^*_{lpn}\\
&\qquad \left. + d_{a_E}^{12}\tilde{\lambda}^{ikj}\tilde{\mu}^*_k/|\mu|\right)\\
(a_\lambda)^{ijk} &= A\,\varepsilon_L\left(c_{A_E}\tilde{\lambda}^{ijk} + \ldots\right)\\
(a_{\lambda'})^{ijk} &= A\,\varepsilon_L\left(c_{A_D}(\tilde{\lambda}')^{ijk} + \ldots\right)\\
a_{\lambda''}^{ijk} &= A\,\varepsilon_B\left(c_{A_{\lambda''}}(\tilde{\lambda}'')^{ijk} + \ldots\right),
\end{aligned}
\tag{54}
$$

up to two flavor spurions and only one in ε_L or ε_B. Just in the case of a_E we consider the expansion up to three spurions.

C Review of RPV contributions to neutrino masses

The neutrino mass matrix receives contributions both at the tree level and from loops [12]. We briefly review for convenience here the general formulae.

C.1 Tree level

In the basis where only h_d develops an electroweak VEV, the tree level contribution to neutrino masses is due to the RPV mixings among neutrinos and higgsinos, proportional to the parameters μ_i. The tree level neutral fermion mass matrix in the basis $(L_\alpha, H_u, \tilde{B}, \tilde{W})$

reads

$$M_\nu = \begin{pmatrix} 0 & m_{RPV} \\ m_{RPV}^T & M_N \end{pmatrix} , \tag{55}$$

where M_N is the 4×4 neutralino mass matrix

$$M_N = \begin{pmatrix} 0 & -\mu & \sin\beta\sin\theta_W M_Z & -\sin\beta\cos\theta_W M_Z \\ -\mu & 0 & -\cos\beta\sin\theta_W M_Z & \cos\beta\cos\theta_W M_Z \\ \sin\beta\sin\theta_W M_Z & -\cos\beta\sin\theta_W M_Z & M_1 & 0 \\ -\sin\beta\cos\theta_W M_Z & \cos\beta\cos\theta_W M_Z & 0 & M_2 \end{pmatrix} \tag{56}$$

and

$$m_{RPV} = \begin{pmatrix} 0 & -\mu_1 & 0 & 0 \\ 0 & -\mu_2 & 0 & 0 \\ 0 & -\mu_3 & 0 & 0 \end{pmatrix} . \tag{57}$$

Under the hypothesis $m_{RPV} \ll M_N$ the matrix M_ν can be perturbatively diagonalized, thus yielding for the three lightest neutrino mass matrix

$$(m_\nu^{(\mathrm{tree})})^{ij} \approx -(m_{RPV} M_N^{-1} m_{RPV}^T)^{ij} = m_\nu^{(\mathrm{tree})} \frac{\mu^i \mu^j}{\sum_i |\mu_i^2|} , \tag{58}$$

where

$$m_\nu^{(\mathrm{tree})} = \frac{\left(M_1\cos^2\theta_W + M_2\sin^2\theta_W\right) M_Z^2\cos^2\beta}{\mu\left((M_1\cos^2\theta_W + M_2\sin^2\theta_W) M_Z^2 \sin 2\beta - M_1 M_2\,\mu\right)} \times \sum_i |\mu_i^2| . \tag{59}$$

By diagonalizing the rank-1 matrix in Eq. (58) we get

$$m_3^{(\mathrm{tree})} = m_\nu^{(\mathrm{tree})} , \qquad m_2^{(\mathrm{tree})} = 0 , \qquad m_1^{(\mathrm{tree})} = 0 . \tag{60}$$

We adopt the following convention for the neutrino mass eigenvalues: $m_3 \geq m_2 \geq m_1$.

C.2 Loops

In order to complete the neutrino spectrum one has to go at the loop level. One-loop neutrino masses get contributions from many diagrams involving different combinations of the coupling μ_i, λ' and λ. On the other hand, under reasonable assumptions on the SUSY parameters (see e.g. [27]), one can focus the attention only on the contribution coming from the trilinear terms λ and λ'. In the basis where the down-quark and the charged-lepton mass matrices are diagonal, one finds [12]

$$(m_\nu^{(\lambda'\lambda')})^{ij} = \frac{3}{16\pi^2} \sum_{k,l,m} \lambda'^{ikl} \lambda'^{jmk}\, \hat{m}_{D_k}\, \frac{(\tilde{m}^{d\,2}_{LR})_{ml}}{m^2_{\tilde{d}_{Rl}} - m^2_{\tilde{d}_{Lm}}} \ln\left(\frac{m^2_{\tilde{d}_{Rl}}}{m^2_{\tilde{d}_{Lm}}}\right) + (i \leftrightarrow j) , \tag{61}$$

$$(m_\nu^{(\lambda\lambda)})^{ij} = \frac{1}{16\pi^2} \sum_{k,l,m} \lambda^{ikl} \lambda^{jmk}\, \hat{m}_{E_k}\, \frac{(\tilde{m}^{e\,2}_{LR})_{ml}}{m^2_{\tilde{e}_{Rl}} - m^2_{\tilde{e}_{Lm}}} \ln\left(\frac{m^2_{\tilde{e}_{Rl}}}{m^2_{\tilde{e}_{Lm}}}\right) + (i \leftrightarrow j) , \tag{62}$$

which at the leading order in the MFV expansion read

$$(m_\nu^{(\lambda'\lambda')})^{ij} = \frac{3}{8\pi^2}\frac{\tilde{m}\,c_{A_D}-\mu\tan\beta}{\tilde{m}^2}\frac{1}{c_{d^c}-c_q}\ln\left(\frac{c_{d^c}}{c_q}\right)(\lambda')^{ikl}(\lambda')^{jlk}\hat{m}_{D_k}\hat{m}_{D_l}\,, \tag{63}$$

$$(m_\nu^{(\lambda\lambda)})^{ij} = \frac{1}{8\pi^2}\frac{\tilde{m}\,c_{A_E}-\mu\tan\beta}{\tilde{m}^2}\frac{1}{c_{e^c}-c_L}\ln\left(\frac{c_{e^c}}{c_L}\right)\lambda^{ikl}\lambda^{jlk}\hat{m}_{E_k}\hat{m}_{E_l}\,. \tag{64}$$

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