Neutrino textures and charged lepton flavour violation in light of $\theta^{13}$, MEG and LHC data

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ABSTRACT

In light of recent results from the LHC, MEG and neutrino experiments, we revisit the issue of charged lepton flavour violation (LFV) in supersymmetric theories with massive neutrinos, where flavour-violating soft supersymmetry-breaking masses for sleptons are induced naturally by radiative corrections. We link our results to the expectations for light neutrinos with a normal mass hierarchy in SU(5), enhanced by an abelian flavour symmetry, with particular focus on $\theta^{13}$. We focus on the radiative decays $\ell_i \rightarrow \ell_j \gamma$ and on detection prospects at the LHC and a linear collider (LC). We use supersymmetric parameters consistent with cosmological considerations and with LHC searches for supersymmetry and the Higgs mass. We find a class of scenarios where the LHC may be sensitive to LFV sparticle decays and LFV processes could be detectable at a LC with centre-of-mass energy above 1 TeV, whereas LFV lepton decays may be suppressed by cancellations in the decay amplitudes.

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1 Introduction

In recent years, the existence of neutrino masses and oscillations with near-maximal $\nu_\mu - \nu_\tau$ and large $\nu_e \to \nu_\mu$ mixing has been established by extensive input from atmospheric [1], solar [2] and long-baseline reactor [3] and accelerator [4, 5] neutrino experiments. Initial input on the possible range of the $\theta_{13}$ [6, 7] was provided by the T2K [8] and MINOS collaborations [9], and definitive evidence for a non-zero value of $\theta_{13}$ has been provided by the reactor experiments Daya Bay [10] and RENO [11], and very recently also by Double Chooz [12].

A natural expectation in theories with massive neutrinos is charged-lepton-flavour violation (LFV), which is enhanced in supersymmetric theories via the renormalization of soft supersymmetry-breaking parameters. The link between neutrino oscillations and violations of the individual lepton numbers $L_{e,\mu,\tau}$ raises the prospect of observing processes such as $\mu \to e\gamma$, $\mu \to 3e$, $\tau \to \mu\gamma$ and $\mu \to e$ conversion on heavy nuclei [13]. The present experimental upper limits on the most interesting of these processes, summarised below, already constrain significantly the parameter spaces of theoretical models:

\begin{align*}
BR(\mu \to e\gamma) &< 5.6 \times 10^{-13} \quad [14], \\
BR(\tau \to \mu\gamma) &< 4.4 \times 10^{-8} \quad [16], \\
BR(\tau \to e\gamma) &< 3.3 \times 10^{-8} \quad [16].
\end{align*}

The strongest constraint on radiative decays is the recent MEG upper limit on $BR(\mu \to e\gamma)$ [14], four times more stringent limit than the previous one [15].

Within the supersymmetric framework, one should also keep in mind other possibilities for observing LFV processes, such as slepton pair production at a Linear Collider (LC) [17, 18, 19, 20, 21, 22, 23, 24] and signals at the LHC [25, 26, 27, 28, 29, 30, 31, 32], particularly in $\chi_2 \rightarrow \chi + e^\pm \mu^\mp$, $\chi_2 \rightarrow \chi + \mu^\pm \tau^\mp$ decays (here $\chi$ is the lightest neutralino, assumed to be the lightest supersymmetric particle (LSP), and $\chi_2$ is the second-lightest neutralino). These decays could provide search prospects that are complementary to direct searches for flavour-violating decays of charged leptons, particularly for heavy superparticle spectra.

In this paper we re-evaluate the prospects for observable charged LFV, based on updated knowledge of neutrino mass and mixing parameters that includes the recent measurement of $\theta_{13}$. We work within the framework of the most natural mechanism for obtaining hierarchical light neutrino masses, namely the see-saw mechanism [33], in which an effective Majorana mass matrix for light neutrinos, $m_{\nu}^{\text{eff}} = m_D^\nu \cdot (M_N)^{-1} \cdot m_D^{\text{DP}}$, arises from Dirac neutrino masses $m_D^\nu$ of the same order as the charged-lepton and quark masses, and heavy Majorana masses $M_N$. In supersymmetric theories, the neutrino Dirac couplings $Y_\nu$ renormalise the soft supersymmetry-breaking sneutrino and slepton masses, generating LFV in a natural way [34]. Even if the soft scalar masses were universal at the unification scale, quantum corrections between the GUT scale and low energies would modify this structure via renormalization-group running, which generates off-diagonal contributions. This effect is particularly interesting in see-saw models, where in general the Dirac neutrino Yukawa couplings cannot be diagonalized simultaneously with the charged-lepton and slepton mass matrices [34]. Given the large mixing of the cor-
responding neutrino species, charged LFV may occur at enhanced rates in supersymmetric extensions of the standard model, giving rise to observable LFV signals \[13, 35, 36, 37, 38, 39\].

We analyse this possibility within the constrained minimal supersymmetric standard model (CMSSM) with universal scalar, gaugino masses and trilinear terms at the GUT scale \(m_0, M_{1/2}\) and \(A_0\), respectively), using mass matrices that are inspired by GUT models with abelian flavour symmetries \[40, 41\]. These textures reproduce naturally the observed fermion mass hierarchies and mixing angles and may also have interesting implications for leptogenesis \[42, 43, 44, 45\]. Despite their phenomenological appeal, however, there are ambiguities and limitations due to the fact that the entries in the mass matrices are determined only up to \(\mathcal{O}(1)\) numerical factors.

The paper is organized in the following way: In Section 2 we look at the theoretical and phenomenological predictions for neutrino mass matrices. In Section 3.1 we discuss the origin of LFV in representative supersymmetric scenarios, in Section 3.2 we analyse the connection between the U(1) charges and LFV and in Section 3.3 we discuss the numerical procedure and the renormalization group runs. The obtained mixing matrices are then used to study various LFV processes: in Section 3.4 we discuss radiative decays and the impact of the new \(\mu \rightarrow e\gamma\) MEG bound on the parameter space of interest; in Section 3.5 we study LFV in \(\chi_2\) decays at the LHC, while LFV from slepton production and decay at a future LC is discussed in Section 3.6. In Section 4 we discuss possible implications for leptogenesis, and finally in Section 5 we summarize the main results of the paper.

## 2 Neutrino Mass Textures inspired by SU(5)

Over the recent years, a plethora of textures have been proposed to explain the data on neutrino masses and mixing. The new data on \(\theta_{13}\) provide additional constraints, excluding certain possibilities and constraining others. Rather than reviewing the vast literature on the subject, we choose a representative model that fits the fermion data and is well-motivated on theoretical grounds. Nevertheless, we try to keep the results as generic as possible, placing emphasis on the links between physical observables. We also keep in mind that several \textit{a priori} different theoretical models may converge to similar phenomenology, since they are matched to the same data.

The example we choose is provided by a SU(5) GUT combined with family symmetries \[40, 41\]. The mass matrices are constructed by looking at the field content of the SU(5) representations, namely: three families of \((Q, u_c, e_c)_i \in 10\) representations, three families of \((L, d_c)_i \in 5\) representations, and heavy right-handed neutrinos in singlet representations. This model therefore has the following properties: (i) the up-quark mass matrix is symmetric, and (ii) the charged-lepton mass matrix is the transpose of the down-quark mass matrix, which relates the mixing of the left-handed leptons to that of the right-handed down-type quarks. Since the observed Cabibbo-Kobayashi-Maskawa (CKM) mixing in the quark sector is due to a mismatch between the mixing of the left-handed up- and down-type quarks, it can be easily reconciled with a large atmospheric neutrino mixing.
Within this framework, and following for example [40, 41, 46], the Yukawa matrices have the form
\[
Y_u \propto \begin{pmatrix} \varepsilon^6 & \varepsilon^5 & \varepsilon^3 \\ \varepsilon^5 & \varepsilon^4 & \varepsilon^2 \\ \varepsilon^3 & \varepsilon^2 & 1 \end{pmatrix}, \quad Y_\ell \propto Y_\ell^T \propto \begin{pmatrix} \varepsilon^4 & \varepsilon^3 & \varepsilon \\ \varepsilon^3 & \varepsilon^2 & 1 \end{pmatrix}, \quad Y_\nu \propto \begin{pmatrix} \varepsilon^{[1\pm n_1]} & \varepsilon^{[1\pm n_2]} & \varepsilon^{[1\pm n_3]} \\ \varepsilon^{[n_1]} & \varepsilon^{[n_2]} & \varepsilon^{[n_3]} \end{pmatrix},
\]
where \(Y_{u,d,\ell,\nu}\) stand for the Yukawa couplings of quarks, charged leptons and neutrinos respectively, and \(n_i\) denote the U(1) charges of the heavy Majorana neutrinos. The heavy Majorana mass matrix is then given by
\[
M_N \propto \begin{pmatrix} \varepsilon^{[2n_1]} & \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_1+n_3]} \\ \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_2]} & \varepsilon^{[n_2+n_3]} \\ \varepsilon^{[n_1+n_3]} & \varepsilon^{[n_2+n_3]} & \varepsilon^{[2n_3]} \end{pmatrix}.
\]

There is no unique choice for the right-handed neutrino charges \(n_1, n_2, n_3\), and several possibilities may be compatible with the low-energy neutrino data. We know, however, that the neutrino masses and mixing angles are related to the \(\nu_L \nu_L^c\) contributions in the effective neutrino mass matrix
\[
m_{\text{eff}} \approx m_{\nu} \frac{1}{M_N} m_{\nu}^D T,
\]
which, if calculated from the matrices in (4) and (5), is of the form:
\[
m_{\text{eff}} \propto \begin{pmatrix} \varepsilon^2 & \varepsilon & \varepsilon \\ \varepsilon & 1 & 1 \\ \varepsilon & 1 & 1 \end{pmatrix}.
\]

This form of \(m_{\text{eff}}\) is quite natural in the simplest see-saw models with a single expansion parameter and generic structures for the heavy and light Majorana mass matrices, due to cancellations that eliminate the dependences on the right-handed charges. Its predictions have been extensively analysed from a phenomenological point of view [47, 48, 49, 50, 51, 52, 53, 54, 55, 56], and give a reasonable match to the data, provided there are no cancellations of potentially large mixing in the charged lepton sector. Among other predictions, \(\theta_{13}\) turns out to be of the correct order of magnitude. It is interesting to also note that, to lowest order in \(\varepsilon\), \(Y_\ell Y_\ell^\dagger\) has the same structure as \(m_{\text{eff}}\), namely
\[
Y_\ell Y_\ell^\dagger \propto m_{\text{eff}} \propto \begin{pmatrix} \varepsilon^2 & \varepsilon & \varepsilon \\ \varepsilon & 1 & 1 \\ \varepsilon & 1 & 1 \end{pmatrix}.
\]

The flavour mixing matrices are determined by the following diagonalizations of the Dirac and Majorana mass matrices:
\[
V_\ell^T (Y_\ell Y_\ell^\dagger) V_\ell^* = \text{diag}(y_\ell^2, y_\mu^2, y_\tau^2),
\]
\[
V_\nu^T (Y_\nu Y_\nu^\dagger) V_\nu^* = \text{diag}(y_\nu^2, y_{\nu_2}^2, y_{\nu_3}^2),
\]
\[
U_N^T M_N U_N = \text{diag}(M_1, M_2, M_3),
\]
\[
U_\nu^T m_{\text{eff}} U_\nu = \text{diag}(m_{\nu_1}, m_{\nu_2}, m_{\nu_3}).
\]
In terms of the above matrices, the Maki-Nakagawa-Sakata (MNS) matrix is given by

\[ U_{MNS} \equiv U = V_c^\dagger U_\nu, \]  

and can be parametrized as:

\[ U = V \cdot \text{diag}(e^{-i\phi_1/2}, e^{-i\phi_2/2}, 1), \]

where

\[ V = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13} \end{pmatrix}, \]

and \( c_{ij} \) and \( s_{ij} \) stand for \( \cos \theta_{ij} \) and \( \sin \theta_{ij} \), respectively.

### 2.1 Predictions for Neutrino Observables

Within the above framework there are ambiguities in the choices of coefficients, limited to a certain extent by requiring consistency with the experimental data. The match of the neutrino data to textures predicted by abelian flavour symmetries can be made by treating coefficients as random variables in Monte-Carlo scans of the multi-dimensional parameter space, in a statistical study of the probability that the textures can naturally reproduced the measured angles and masses \[49, 50, 51, 52, 53, 54, 55\].

Here, being interested in matrices that are naturally consistent with the neutrino data, we proceed by taking the expansion parameter to be \( \varepsilon = 0.2 \) and multiply the entries of \( Y_\ell, Y_\nu \) and \( M_N \) in Eqs. (4) and (5) by coefficients \( \ell_{ij}, \nu_{ij} \) and \( N_{ij} \) in the range \( \pm [0.5, 2] \). In addition, we impose the following constraints:

(i) We select the charged lepton Yukawa coupling matrix \( Y_\ell \) so that the correct charged-lepton mass hierarchies are reproduced, namely

\[ \frac{m_\mu}{m_\tau} \sim 0.06, \quad \frac{m_e}{m_\tau} \sim 2.5 \times 10^{-3}. \]

(ii) \( Y_\nu \) and \( M_N \) are required to give a light neutrino mass matrix \( m_{\text{eff}} \) of the form (7), with entries that deviate by a factor \( \varepsilon \in [0.5, 2] \) from those in Eq. (7).

(iii) We impose normal hierarchy among the neutrino masses. We fix \( m_{\nu_3} \sim \sqrt{\Delta m_{\text{atm}}^2} \sim 0.05 \) eV and require \( 0.16 < m_{\nu_2}/m_{\nu_3} < 0.19, m_{\nu_1} < 0.2m_{\nu_2} \) consistent with the measured values of \( \Delta m_{\text{sol}}^2 \) and \( \Delta m_{\text{atm}}^2 \) \[57\].

(iv) We require the following range of mixing angles \[7\]:

\[ 0.27 < \sin^2 \theta_{12} < 0.35, \]
\[ 0.34 < \sin^2 \theta_{23} < 0.67, \]
\[ 0.018 < \sin^2 \theta_{13} < 0.033. \]

The range on \( \theta_{13} \) is consistent with the values reported by both \[10, 11\] at the 3\( \sigma \) level.
the CP-violating phase $\delta$ and $\theta$ predictions without the RGE runs).

REAP using the code into account renormalization effects. The RGE runs for the “see-saw” MSSM are evaluated in all cases, the effect of varying $\tan\beta$ has an impact of less than 2% on the final value of the neutrino mixing angles.

Table 1: Indicative textures for $Y_t$, $Y_\nu$ and $M_N$ at the GUT scale, to be studied in detail below. The $n_i$ are abelian charges, that can only be constrained by LFV. The computation of the neutrino mixing angles includes the RGE effects using $\tan\beta = 45$ (in parenthesis we quote the predictions without the RGE runs).

(v) We make a further selection by requiring that the hierarchy of eigenvalues of $Y_t^r Y_t$ (which, as discussed above, has a similar structure to $m_{eff}$ and $Y_\nu Y_\nu^r$) preserves the order of the gauge eigenstates. This reduces the density of solutions in the plots and implies large off-diagonal elements in both $V_t$ and $U_\nu$.

The selection of coefficients in the textures is performed so that the above conditions are satisfied and the coefficients are chosen to be real in the range $\pm[0.5, 2]$. Given our ignorance of the CP-violating phase $\delta$, we focus on the case $\delta = 0$ and do not include the Majorana phases $\phi_{1,2}$.

In Table 1 we provide two representative examples of our fits, which will be used for our analysis below. We quote our predictions for neutrino mixing angles with and without taking into account renormalization effects. The RGE runs for the “see-saw” MSSM are evaluated using the code REAP, described in Ref. [69]. The coefficients are taken at the GUT scale. We work with $\tan\beta = 45$, since this is the largest value that we will use in the numerical computations of the next section (and larger $\tan\beta = 45$ results to larger corrections). In all cases, the effect of varying $\tan\beta$ in the range of values used in our examples (from 16 to 45) has an impact of less than 2% on the final value of the neutrino mixing angles.

In Fig. 1 we present the predictions for the neutrino mixing angles corresponding to the above criteria. Here we do not take into account the RGE run of the mass matrices, since these effects will not affect the global picture of the solutions displayed in Fig. 1. Within this class of models, most of the solutions that reproduce the correct range of $\theta_{12}$ and $\theta_{23}$, also predict a neutrino mixing angle $\theta_{13}$ that is compatible with the data from [10,11]. Within the range of $\theta_{13}$ the model predictions are mostly in the mid-lower range of $\theta_{12}$. On the other hand, in the case of $\theta_{23}$, a higher density of solutions is found in the mid-higher range of $\theta_{23}$. We note that the number of fits predicting the maximal value $\theta_{23} = \pi/4$ is smaller when we impose the hierarchy

| Fit | $Y_t$ | $Y_\nu$ | $M_N$ |
|-----|-------|--------|-------|
| 1   | $\begin{pmatrix} \varepsilon^4 & -1.6\varepsilon^3 & 1.2\varepsilon \\ 0.7\varepsilon^3 & 1.6\varepsilon^2 & -0.6 \\ 0.7\varepsilon^3 & -1.7\varepsilon^2 & -1.3 \end{pmatrix}$ | $\begin{pmatrix} \varepsilon^{[1+n_1]} & \varepsilon^{[1+n_2]} & -1\varepsilon^{[1+n_3]} \\ 0.8\varepsilon^{[n_1]} & \varepsilon^{[n_2]} & -1.2\varepsilon^{[n_3]} \\ -1.3\varepsilon^{[n_1]} & \varepsilon^{[n_2]} & 0.7\varepsilon^{[n_3]} \end{pmatrix}$ | $\begin{pmatrix} \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_1+n_3]} & -1.7\varepsilon^{[n_1+n_3]} \\ \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_1+n_3]} & \varepsilon^{[n_2+n_3]} \\ -1.7\varepsilon^{[n_1+n_3]} & \varepsilon^{[n_2+n_3]} & -2\varepsilon^{[n_3]} \end{pmatrix}$ |
| 2   | $\begin{pmatrix} \varepsilon^4 & -1.5\varepsilon^3 & -2\varepsilon \\ \varepsilon^3 & -1.9\varepsilon^2 & 0.5 \\ 0.5\varepsilon^3 & -\varepsilon^2 & 0.75 \end{pmatrix}$ | $\begin{pmatrix} \varepsilon^{[1+n_1]} & \varepsilon^{[1+n_2]} & -2\varepsilon^{[1+n_3]} \\ 1.5\varepsilon^{[n_1]} & \varepsilon^{[n_2]} & -0.75\varepsilon^{[n_3]} \\ 1.9\varepsilon^{[n_1]} & \varepsilon^{[n_2]} & 1.5\varepsilon^{[n_3]} \end{pmatrix}$ | $\begin{pmatrix} \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_1+n_3]} & -1.9\varepsilon^{[n_1+n_3]} \\ \varepsilon^{[n_1+n_2]} & \varepsilon^{[n_1+n_3]} & \varepsilon^{[n_2+n_3]} \\ -1.9\varepsilon^{[n_1+n_3]} & \varepsilon^{[n_2+n_3]} & 1.9\varepsilon^{[n_3]} \end{pmatrix}$ |

$\sin^2\theta_{13}=0.020(0.022)$, $\sin^2\theta_{12}=0.267(0.274)$, $\sin^2\theta_{23}=0.580(0.613)$
Figure 1: *We summarise the correlations between the neutrino mixing angles before and after constraining the model coefficients (as discussed in the text). The solid lines indicate the experimental bounds, and the small black crosses represent models satisfying all constraints. The two large magenta crosses correspond to the benchmark models that are discussed in the text and in Table 1 and are used for numerical calculations.*

condition \((iii)\), as compared to the case where only the experimental bounds on \(\Delta^2 m_{\text{sol}}\) and \(\Delta^2 m_{\text{atm}}\) are considered. This is consistent with the observed deviation of \(\theta_{23}\) from its maximal value \([57]\).

As benchmarks for studying LFV in subsequent sections, we have identified the two sets of textures of Table 1 and indicated with crosses in Fig. 1. These benchmarks are chosen as representatives of the two different regions identified in the global statistical data analysis performed in \([57]\) and also shown in Fig 1. Fit 1 lies in the right region of the right panel of Fig 1 with larger \(\theta_{23}\), and Fit 2 lies in the left region of the same panel, with smaller \(\theta_{23}\).

We would like to point out that the fits in Fig. 1 are independent of the charges \(n_i\), which affect \(Y_\nu\) and \(M_N\) but not their combination in \(m_{\text{eff}}\). On the other hand, the choices of \(n_i\) do affect the rates for charged LFV processes, as we will show in the next Section. Charged LFV processes are therefore powerful probes of parameters that cannot be constrained by lepton hierarchies and better measurements of the neutrino parameters.
3 Charged-Lepton-Flavour Violation

3.1 Slepton masses in see-saw models.

The neutrino mass textures discussed above arise naturally from $SU(5)$ enhanced by a U(1) flavour symmetry. In order to study LFV processes, however, we are guided to a large extent by phenomenological considerations and thus our results are more generic. The benchmark solutions previously identified, which naturally reproduce the correct neutrino phenomenology, are used to obtain the matrices that determine the LFV vertices in the context of the CMSSM (extended with right-handed heavy Majorana neutrinos arising from a see-saw mechanism).

Since the Dirac neutrino and charged-lepton Yukawa couplings cannot, in general, be diagonalized simultaneously (and since both types of lepton Yukawa couplings appear in the RGEs) the lepton Yukawa matrices and the slepton mass matrices at low energies cannot be diagonalized simultaneously, either. In the basis where the charged-lepton masses $m_\ell$ are diagonal, the soft slepton-mass matrix acquires corrections that contain off-diagonal contributions from renormalization at scales below $M_{GUT}$, of the following form in the leading-log approximation [35]:

$$\delta m_\\tilde{\ell}^2 \propto \frac{1}{16\pi^2}(6m_0^2 + 2A_0^2)Y_\nu^\dagger Y_\nu \ln\left(\frac{M_{GUT}}{M_N}\right), \quad (18)$$

Here $M_N$ is the intermediate scale where the effective neutrino-mass operator is formed. The physical charged slepton masses are obtained by numerical diagonalization of the following $6 \times 6$ matrix:

$$m_\\tilde{\ell}^2 = \begin{pmatrix} m_{LL}^2 & m_{LR}^2 \\ m_{RL}^2 & m_{RR}^2 \end{pmatrix}, \quad (19)$$

where all the entries are $3 \times 3$ matrices in flavour space. Using the basis where $Y_\ell$ is diagonal, it is convenient to write the $3 \times 3$ entries of (19) in the form:

$$m_{LL}^2 = (m_\\tilde{\ell}^\text{diag})^2 + \delta m_\tilde{\ell}^2 + m_\ell^2 - \frac{1}{2}(2M_W^2 - M_Z^2)\cos 2\beta,$$  \quad (20)

$$m_{RR}^2 = (m_\\tilde{\ell}_R^\text{diag})^2 + m_\ell^2 - (M_Z^2 - M_W^2)\cos 2\beta,$$  \quad (21)

$$m_{RL}^2 = (A_\ell^\text{diag} + \delta A_\ell - \mu \tan \beta)m_\ell,$$  \quad (22)

$$m_{LR}^2 = m_{RL}^2.$$  \quad (23)

Here $\tan \beta$ is the ratio of the two MSSM Higgs vevs, $(m_\\tilde{\ell}^\text{diag})^2, (m_\\tilde{\ell}_R^\text{diag})^2$ and $A_\ell^\text{diag}$ denote the diagonal contributions to the corresponding matrices, obtained by numerical integration of the RGEs, and $\delta m_\tilde{\ell}^2$ and $\delta A_\ell$ denote the corrections to off-diagonal terms that appear because $Y_\nu$ and $Y_\ell$ cannot be diagonalized simultaneously.

The full mass matrix for left- and right-handed sneutrinos has a $12 \times 12$ structure, given in terms of $3 \times 3$ Dirac, Majorana and sneutrino mass matrices. The effective $3 \times 3$ mass-squared matrix for the left-handed sneutrinos has the same form as the $m_{LL}^2$ part (23) of the $6 \times 6$ charged-slepton matrix (19), with the difference that now the Dirac masses are absent. In
Ref. [36] it was shown that is sufficient to use

$$m_ν^2 = (m_ν^{\text{diag}})^2 + δ m_ν^2 + \frac{1}{2} M_2^2 \cos 2β .$$

The matrix responsible for LFV in the lepton-slepton-gaugino vertices is

$$V_{LFV} = V_D^\dagger V_ℓ$$

and the slepton mass matrices contain off-diagonal terms generated by:

$$m_{LL}^2 = V_{LFV}^\dagger (m_{LL}^2)_{d} V_{LFV} ,$$

while the $A$-terms become:

$$A_ℓ = V_{LFV}^T (A_ℓ)_{d} .$$

Here $(m_{LL}^2)_{d}$ and $(A_ℓ)_{d}$ are the terms resulting from the RGE running of the universal soft terms at the GUT scale in a basis where $Y_ν$ is diagonal. The corresponding effects in $m_{RR}^2$ are negligible and are not considered in the numerical calculations.

We remark that, in general, in the framework of susy SU(5) GUT with U(1) family symmetries, flavour dynamics are linked to scalar singlet fields, flavons, whose non-zero vacuum expectation value breaks the U(1) symmetry. The RGE running of the parameters above the GUT scale due to flavons dynamic induces flavour dependent corrections to sfermion soft mass matrices and A-terms and thus potentially large LFV effects [36, 58, 59, 60, 61]. However, while flavon effects can be potentially very large, we know from flavour phenomenology that this is not the case and that they have to be suppressed to the point that they are comparable to the effects we consider here. Such a suppression can be achieved, among others, in an scenario where the effect of non-universal soft terms is diminished by RGE effects beyond the GUT scale, as indicated in [46]. Given that the exact knowledge of flavon effects depends on model building and physics of unknown scales (supersymmetry breaking scale, flavon dynamics scale or string scale), a complete mechanism cancelling the undesired flavour-violating soft terms at $M_{GUT}$ goes beyond the scope of this paper.

### 3.2 LFV and neutrinos

We now study the conditions under which the favoured range of neutrino masses and mixing, can lead in a natural way to observable signatures for charged LFV. As already discussed, while the neutrino parameters are independent of the charges $n_i$, this is not the case for LFV. As a result, LFV can provide a way to probe the right-handed neutrino sector, for which only limited information is available.

In section 2.1, we identified two representative benchmark fits suitable for studying charged LFV. The level of charged LFV is determined by the product $V_{LFV} = V_D^\dagger V_ℓ$, and thus by the charges $n_i$ which enter in the Dirac neutrino mixing matrix $V_D^\dagger$. Having a $V_ℓ$ with large off-diagonal 1-2 and 2-3 entries is a natural choice to match the lepton data. Then, different
choices of \( n_i \) lead to different possibilities for \( V_D^\dagger \); in fact, there are two possibilities associated with a \( V_{\ell} \) with large off-diagonal elements: (1) the charge combinations generate a \( V_D^\dagger \) with small off-diagonal elements; in this case, \( V_{LFV} \sim V_{\ell} \). (2) the off-diagonal elements of \( V_D^\dagger \) are large, but multiplied with \( V_{\ell} \), they can give either large or small elements in \( V_{LFV} \) depending on coefficients and phases.

An illustration of the dependences of the entries in \( V_{LFV} \) on the different right-handed neutrino charges is given in Table 2. For simplicity, we focus on Fit 2, noting that similar results hold for Fit 1. The matrix (i) is an example of case (1), with small off-diagonal elements in \( V_D \). In (ii) and (iii), \( V_D \) has large off-diagonal elements which enhance \( V_{LFV} \). Finally in (iv), \( V_D \) has large off-diagonal elements but cancellations with \( V_\ell \) occur in the 2-3 sector, suppressing LFV.

### Table 2: Values for the matrix \( V_{LFV} \) of Eq. (25) corresponding to Fit 2 of Table 1 with \( \varepsilon = 0.2 \) and different choices of the \( U(1) \) charges, Eqs. (4, 5).

| \( n_i \) | (i) \( \{n_1 = 1, n_2 = 0, n_3 = 0\} \) | (ii) \( \{n_1 = 2, n_2 = 1, n_3 = 0\} \) | (iii) \( \{n_1 = 2, n_2 = 0, n_3 = 1\} \) | (iv) \( \{n_1 = 0, n_2 = 1, n_3 = 0\} \) |
|---|---|---|---|---|
| \( V_{LFV} \) | \( \begin{pmatrix} 0.805 & -0.385 & -0.451 \\ 0.182 & 0.885 & -0.429 \\ 0.565 & 0.263 & 0.782 \end{pmatrix} \) | \( \begin{pmatrix} 0.805 & -0.385 & -0.452 \\ -0.064 & 0.700 & -0.711 \\ 0.590 & 0.601 & 0.539 \end{pmatrix} \) | \( \begin{pmatrix} -0.805 & 0.384 & 0.453 \\ 0.544 & 0.782 & 0.305 \\ -0.237 & 0.492 & -0.838 \end{pmatrix} \) | \( \begin{pmatrix} 0.806 & -0.401 & -0.436 \\ -0.437 & -0.899 & 0.016 \\ -0.399 & 0.178 & -0.901 \end{pmatrix} \) |

#### 3.3 Numerical procedure and RGEs

The recent LHC measurement of the Higgs mass \([63, 64]\) imposes severe constraints in the CMSSM parameter space. More specifically, a Higgs masses of \( m_h \sim 125 \) GeV implies, in general, a relatively heavy sparticle spectrum, which is consistent with the cosmological constraint on the neutralino relic density only in limited regions. A global analysis of the CMSSM parameter space was performed in [62], yielding two almost equally good fits to the available data, one with relatively low sparticle masses and \( \tan \beta \sim 16 \), and the other with larger sparticle masses and \( \tan \beta \sim 45 \).

\[
\begin{align*}
(a) & \quad \tan \beta = 16, \quad m_0 = 300 \text{ GeV}, \quad M_{1/2} = 910 \text{ GeV}, \quad A_0 = 1320 \text{ GeV}, \\
(b) & \quad \tan \beta = 45, \quad m_0 = 1070 \text{ GeV}, \quad M_{1/2} = 1890 \text{ GeV}, \quad A_0 = 1020 \text{ GeV}.
\end{align*}
\]

The sign of \( \mu \) is positive, as favored by \( g_\mu - 2 \) measurements. Regarding cosmological considerations, point (a) belongs to the area where the WMAP-favoured range of \( \Omega \chi h^2 \) is obtained via \( \chi - \tilde{\tau} \) coannihilation\(^2\) whereas point (b) lies in the funnel region where the neutralino LSP annihilates rapidly via direct-channel \( H/A \) poles.

\(^1\)Note that our \( A_0 \) values have opposite sign with respect to those of Ref. [62], where the authors use a definition for the trilinear scalar coupling that differs from the one in standard codes like Suspect and SoftSusy.

\(^2\)We note that in this region the \( \chi - \tilde{\tau} \) mass difference is very small, offering other experimental challenges and opportunities \([65, 66]\).
We evaluate the RGEs using universal soft terms at the GUT scale, \( M_{\text{GUT}} \). The standard model parameters are evaluated at \( M_Z \) and \( m_t(m_t) \). At the GUT scale, defined as the meeting point of the gauge couplings \( g_1 \) and \( g_2 \) (\( g_3 \) is set so that \( \alpha_s(M_Z) = 0.1172 \)), we work in a basis where \( Y_\nu \) is diagonal. Non-diagonal elements of the soft mass matrices are induced from the fact that \( Y_\ell \) cannot be diagonalized simultaneously with \( Y_\nu \). The right handed neutrino scale is identified with the mass of the largest eigenvalue of \( M_N \), \( M_3 \). The coupling \( Y_{\nu_3} \) is calculated by requiring that \( m_{\nu_3} = 0.05 \text{ eV} \) at low energy, using the respective RGE \([67, 38, 68]\). At \( M_3 \) we decouple the see-saw parameters from the RGE; in doing so, we neglect the effect of the lighter neutrinos, which in the case of hierarchical neutrinos is not large (even if \( M_2 \) and \( M_3 \) are much lighter than \( M_3 \), the corresponding \( Y_\nu \) must decrease according to the see-saw relation, resulting to an insignificant impact on the slepton mass running). At \( M_3 \) we rotate all the fields in the basis where \( Y_\ell \) becomes diagonal; in this basis, \( m_L^2 \) and \( A_\ell \) take the form of Eqs. (26) and (27) while \( m_{RR}^2 \) remains essentially diagonal since its RGE is not affected by \( Y_\nu \). Moreover, in this basis, only the diagonal terms evolve from \( M_3 \) down to low energies. The matrix \( V_{\text{LFV}} \) is computed using Yukawa textures that match the neutrino data. In our RGE analysis we do take into account the change of the overall scale of \( m_{\text{sfit}} \) and \( Y_\ell \) but not the RGE dependence of each matrix element (for hierarchical neutrinos this dependence is small and can be absorbed in the uncertainty of the coefficients used to fit the texture without a significant effect in the slepton mass matrices). At low energies, we decouple the SUSY particles at \( M_{\text{SUSY}} = \sqrt{m_\tilde{\tau}_1 \cdot m_\tilde{\tau}_2} \) and continue with the SM RGEs to \( M_t \) and \( M_Z \), with the initial conditions of the RGEs being iteratively adjusted to the experimental data.

### 3.4 Predictions for Radiative Decays

The matrix element of the electromagnetic-current operator between two distinct lepton mass eigenstates \( \ell_i \) and \( \ell_j \) is given in general by

\[
\mathcal{T}_\lambda = \langle \ell_i | (p - q) | J_\lambda | \ell_j (p) \rangle = \bar{u}_i (p - q) m_j \sigma_{\lambda \beta} q^\beta (A_M^L P_L + A_M^R P_R) + (q^2 \gamma_\lambda - q \gamma \cdot q) (A_E^L P_L + A_E^R P_R) u_j (p),
\]

where \( q \) is the photon momentum. The coefficients \( A_M \) and \( A_E \) denote contributions from neutralino/charged-slepton and chargino/sneutrino exchanges. The amplitude of the LFV process is then proportional to \( \mathcal{T}_\lambda \epsilon^\lambda \), where \( \epsilon^\lambda \) is the photon-polarization vector. The branching ratios (BR) of the decays \( \ell_j \rightarrow \ell_i + \gamma \) are calculated using the exact formulas of Ref. [35].

In Fig. 2 we show numerical predictions for the LFV branching ratios arising from the textures introduced above and for the CMSSM parameters specified in (28). We show the effect of varying \( M_3 \) from \( 6 \times 10^{14} \text{ GeV} \) down to \( 10^{12} \text{ GeV} \), for Fit 2 of Table 1 and for the choices of right-handed neutrino charges of Table 2. We can see that the experimental upper bounds on BR(\( \mu \rightarrow e\gamma \)) can be reached with some of the Yukawa textures we studied, even with the heavy sparticle spectrum implied by the benchmark point \((a)\). The new MEG bound on BR(\( \mu \rightarrow e\gamma \)) imposes constraints on the see-saw scale for all charge choices of Table 2 for point \((a)\), and for fit \((ii)\) at point \((b)\).
Figure 2: Predictions for the rare LFV decays $\ell_i \rightarrow \ell_j \gamma$ as a function of the right-handed neutrino mass $M_N$, for the benchmark points displayed in (28) (a) (thick line), (b) (thin line), using the neutrino mixing fits shown in Table 2. The solid lines correspond to case (iv), dashed ones to (iii), dot-dash to (ii) and dot to (i). The horizontal solid lines indicate the current experimental upper bounds, while the dashed line correspond to the previous MEG limit on $BR(\mu \rightarrow e \gamma)$. 
In the case of point (a), we find cancellations that reduce the branching ratios for some neutrino mixing fits. This happens because of the large value of $A_0$, which leads to significant cancellations among the different LFV decay amplitudes $[35, 36, 37]$. For the fits (i) and (ii) we find this type of cancellation in $\text{BR}(\tau \to \mu \gamma)$ and $\text{BR}(\tau \to e \gamma)$, for a range of values of the right-handed neutrino scale $M_3$. This is due to the fact that, at large values of $A_0$, the contribution $A_M^R$ from the neutralino/charged slepton loops in eq. (29) cancels with the one arising from chargino/sneutrino loops, $A_E^R$, which is the dominant contribution for small $A_0$. The ratio of these two contributions can be modulated by the parameters that determine the size of the flavor mixing elements $\delta m^2_{\tilde{\ell}}$ in eq. (23). In our case the scale of $M_3$ also determines the strength of $Y_\nu$ and thus the size of the LFV terms. We use Fig. 2 and the current MEG bound on $\text{BR}(\mu \to e \gamma)$ to fix the $M_3$ scale for further studies: $M_3 = 2 \times 10^{13}$ GeV for the benchmark point (a) and $M_3 = 10^{14}$ GeV for the benchmark point (b).

### 3.5 LFV in $\chi_2$ decays at the LHC

A promising channel to search for LFV at the LHC is the production and decay of the second lightest neutralino, $\chi_2 \to \chi + \tau^\pm + \mu^\mp$. In $[26, 29]$ it was shown that in order to have a signal that could be distinguished from the background, the ratio

$$R_{\tau\mu} = \frac{\Gamma(\chi_2 \to \chi + \tau^\pm + \mu^\mp)}{\Gamma(\chi_2 \to \chi + \tau^\pm + \tau^\mp)}$$

should be of the order of 10%. For $A_0 = 0$, due to the absence of cancellations suppressing rare charged lepton decays, one had to go beyond the CMSSM to find solutions compatible with all experimental and cosmological data $[29]$. Here, we extend this study to large values of $A_0$, noting that the cancellations that can arise in the branching ratios of radiative decays do not occur in $R_{\tau\mu}$. This opens the possibility to observe LFV in neutralino decays at the LHC, in cases where LFV would be undetectable in rare charged lepton radiative decays.

To see whether this is indeed the case, we proceed with the computation including all contributing on-shell sfermion exchange diagrams, as given in $[28]$:

$$\text{BR}(\chi_2 \to \chi \tau^\pm \mu^\mp) = \sum_{i=1}^3 \left[ \text{BR}(\chi_2 \to \tilde{\ell}_i \mu) \text{BR}(\tilde{\ell}_i \to \tau \chi) + \text{BR}(\chi_2 \to \tilde{\ell}_i \tau) \text{BR}(\tilde{\ell}_i \to \mu \chi) \right]$$

These are evaluated in the benchmark points (a) and (b), to see whether the branching ratio can be of the order of the required reference value.

In Fig. 3 we present the predictions for the branching ratio (30) as a function of $M_3$. For point (a), our predictions are within the reach of the LHC for values of $M_3$ that are compatible with the MEG limit. For point (b), the predictions are below the expected experimental sensitivity.

### 3.6 LFV at a Linear Collider

In supersymmetric models where LFV is produced by lepton-slepton vertices, observable signatures may occur either directly, in slepton-pair production, or indirectly, via slepton production
in cascade decays \cite{17}. If the flavour mixing is introduced in the left-left slepton sector, as is the case for the models under consideration here, the dominant channels are slepton-pair production and LFV decays, such as:

\[
\begin{align*}
    e^+ e^- & \rightarrow \tilde{\ell}_i^- \tilde{\ell}_j^+ \rightarrow \tau^+ \mu^- \chi_1^0 \chi_1^0, \\
    e^+ e^- & \rightarrow \nu_i \nu_j^c \rightarrow \tau^+ \mu^- \tilde{\chi}_1^+ \tilde{\chi}_1^- .
\end{align*}
\]  

In the CMSSM benchmark points introduced above, the channel mediated by charged sleptons clearly dominates over the sneutrino-pair production process, and may lead to a cross section of the order of 1 fb; this is the reference value of \cite{23}, for a LFV signal of $\mu^\pm \tau^\pm$ pairs that can be distinguished from the background, according to the study made in \cite{21}. Here, we extend our previous results \cite{23} which were focused on the production of $\mu^\pm \tau^\pm$ pairs by considering the full structure of the Yukawa matrices, thus comparing the LFV production of charged leptons of all generations. Complete expressions for the LFV cross sections are given in Ref. \cite{18} and used in our work.

In Fig. 4 we present the expected cross sections $\sigma(e^+ e^- \rightarrow \tilde{\ell}_i^- \tilde{\ell}_j^+ \rightarrow \ell_a^\pm \ell_b^\mp + 2\chi^0)$ as a function of $\sqrt{s}$ for the same choice of parameters as in Figures 2 and 3. Naturally, the cross sections in the case of point (a) are larger because sleptons and gauginos are much lighter than in the spectrum of point (b). In (a) sleptons and sneutrinos are nearly degenerate and the cross sections, at energies above the threshold for pair production that is around 1.2 TeV, show a feeble decrease with $\sqrt{s}$. Final states with $e\mu$ pairs have the largest cross section, with value between 1 fb and 10 fb, with a small dependence on the choice of charges (i)-(iv). On the other hand, the cross section for the processes with $\tau e$ and $\tau \mu$ final states show a stronger dependence on the choice of charges, varying between $10^{-1}$ fb and 1 fb in the first case, and between $10^{-2}$ and 1 fb in the second case. Similar behaviour is observed in the case of point (b) where the heavy spectrum implies a threshold around 3 TeV and cross sections below $10^{-1}$ fb.

Figure 3: The ratio defined in Eq. (30) is presented for the CMSSM points (a) (thick line) and (b) (thin line) (in Eq. 28), with the same notation as in Fig. 2.
Figure 4: Values of the cross sections $\sigma(e^+e^- \rightarrow \tilde{\ell}_i^+ \tilde{\ell}_j^- \rightarrow \ell_a^\pm \ell_b^\mp + 2 \chi^0)$ ($\ell_a \neq \ell_b$ as indicated in each panel) as functions of $\sqrt{s}$. The line styles are the same as those in Fig. 2. For point (a) we use $M_3 = 2 \times 10^{13}$ GeV, while for the point (b) we work with $M_3 = 10^{14}$ GeV.

According to Fig. 2 at the selected value of $M_3 = 2 \times 10^{13}$ GeV, $\text{BR}(\tau \rightarrow \mu \gamma)$ and $\text{BR}(\tau \rightarrow e \gamma)$ are suppressed. Since these cancellations do not occur for the LFV LC signals, it is possible to observe slepton flavor oscillations at the LC, in cases where LFV would be undetectable in rare charged lepton decays (as it could also happen at the LHC). It is worth to remark that the CLIC project for a linear collider has as nominal centre of mass energies the values 1.4 TeV and 3 TeV [70, 71], with the option of reaching 5 TeV. The value $\sqrt{s} = 1.4$ TeV is optimal for point (a) where the LFV cross sections are nearly maximal.
4 LFV and Leptogenesis

We comment now on possible links between our LFV predictions and leptogenesis [42] through the decays of heavy, right-handed Majorana neutrinos into leptons and antileptons. Since LFV is related to the see-saw parameters in our framework, there can be interesting consequences for LFV in charged lepton decays and elsewhere [43].

In previous Sections, we have used real parameters to fit the Yukawa couplings, but small phases that would not alter our LFV considerations could induce significant contributions to the lepton and baryon asymmetries of the universe. In what follows, we explore what sizes of the phases in $Y_\nu$ can predict a value for the baryon asymmetry $Y_B$ compatible with the observation [72]

$$Y_B = (6.16 \pm 0.16) \times 10^{-10}.$$  

(33)

For hierarchical heavy neutrinos in a supersymmetric see-saw model, one has [44],

$$Y_B \simeq -10^{-2} \kappa \epsilon_1,$$  

(34)

where $\epsilon_1$ is the CP-violating asymmetry in the decay of the lightest Majorana neutrino and $\kappa$ an efficiency factor parametrizing the level of washout of the generated asymmetry by inverse decay and scattering interactions. The latter depends on the mass of the decaying neutrino $M_1$ and the effective mass parameter

$$\tilde{m}_1 = \frac{v_u^2}{M_1} (\lambda^\dagger \lambda)_{11},$$  

(35)

where $\lambda_\nu$ is the Dirac neutrino Yukawa matrix in the basis where the Majorana masses are diagonal, and $v_u$ is the vev of the Higgs field that couples to up-quarks and neutrinos.

The CP-violating decay asymmetry $\epsilon_1$ arises from the interference between tree-level and one-loop amplitudes:

$$\epsilon_1 = \frac{1}{(8\pi \lambda^\dagger \lambda)_{11}} \sum_{i \neq 1} \text{Im} \left[ \left( (\lambda^\dagger \lambda)_{1i} \right)^2 \right] f \left( \frac{M_1^2}{M_i^2} \right),$$  

(36)

with $f(y) = \sqrt{y} \left[ \frac{1}{1-y} + 1 - (1 + y) \ln \left( \frac{1+y}{y} \right) \right]$. The value of the CP asymmetry depends on the details of the model, but a model-independent upper bound exists, given by [45]

$$|\epsilon_1| \leq \frac{3}{8\pi} \frac{M_1}{v_u^2} (m_3 - m_1),$$  

(37)

where the $m_i$ are the masses of the light neutrinos.

In our work, the flavour and GUT symmetries enable us to correlate $M_1$ with the mass of the heaviest right-handed neutrino, $M_3$, which controls the LFV effects. The neutrino Yukawa couplings of all generations are also related. Then, using as a first approximation eq. (37) and $\kappa$ as derived in [44], we can infer how leptogenesis may be accommodated in our study. Our
Some typical results are presented in Table 3, where we see that at a reference value of $M_3 = 5 \cdot 10^{13}$ GeV and the lower to $M_3 = 10^{12}$ GeV.

fits predict large $\tilde{m}_1$ that, according to [14], implies a strong wash-out regime in which $\kappa$ ranges between $\sim 10^{-3}$ and $\sim 10^{-4}$.



| $M_1$(GeV) | (i) | (ii) | (iii) | (iv) |
|------------|-----|------|-------|------|
| $8.6 \cdot 10^{10}$ | $4.3 \cdot 10^{12}$ | $2.6 \cdot 10^{11}$ | $5.4 \cdot 10^{11}$ | $2.3 \cdot 10^{12}$ |
| $m_1$(eV) | 0.19 | 0.78 | 5.17 | 1.19 |
| $Y^\alpha_{max}$ | $1.0 \cdot 10^{-8}$ | $1.2 \cdot 10^{-10}$ | $2.8 \cdot 10^{-11}$ | $6.6 \cdot 10^{-10}$ |
| $Y_B^*$ | $1.3 \cdot 10^{-10}$ | $3.5 \cdot 10^{-11}$ | $1.2 \cdot 10^{-12}$ | $3.2 \cdot 10^{-12}$ |

Table 3: Baryon asymmetry predictions based on four representative fits. Here, $Y^\alpha_{max}$ is the value obtained using eq. (37), and $Y_B^*$ is the prediction for $Y_B$ using eq. (36) and inserting a phase of 0.1 rad in the (12) element of $Y_\nu$. In each row the upper value corresponds to $M_3 = 5 \cdot 10^{13}$ GeV and the lower to $M_3 = 10^{12}$ GeV.

Some typical results are presented in Table 3, where we see that at a reference value of $M_3 = 5 \cdot 10^{13}$ GeV, $Y^\alpha_{max}$ (calculated using $\epsilon_1$ from eq. (37)) is considerably larger than the experimental value of $Y_B$ for fit (i), implying that in this case the CP-violating phases have to be small enough for $\epsilon_1$ to be well below its maximal value. On the other hand, $Y_B^*$ is below $Y_B$ for fit (iii), and of the same order of magnitude for fits (ii) and (iv). These differences between the fits are due to the different hierarchies between the heavy Majorana masses and the neutrino Yukawa couplings in $\tilde{m}_1$ (which are determined by the right-handed neutrino charges). Consequently, they indicate how leptogenesis can be used as an additional probe of the right-handed neutrino sector, for which very limited information is provided by the neutrino data alone.

Looking at the predictions for leptogenesis in more detail and using the complete expression for $\epsilon_1$ in (36), we see that fit (i) can accommodate comfortably the observed baryon asymmetry $Y_B$ with phases of $O(0.1)$ rad, which would not change the LFV predictions. The remaining three models, if the phases are small, would under-produce $Y_B$. We also note that decreasing the scale $M_3$ would decrease both $Y_B^*$ and the LFV effects, whereas increasing $M_3$ to values that would correspond to the perturbative limit for $Y_\nu$ would increase $Y_B^*$, but not sufficiently to reach its experimental value with small phases [3]. In these cases, either one would have to postulate an additional source of baryon asymmetry, or one should explore predictions for LFV in the presence of large phases. Examining these possibilities lies beyond the scope of this paper. However, we do note that overproduction of baryons is not a problem in our scenario, even in the absence of extra sources of entropy.

Furthermore, in the case of benchmark point (a) we can see in Fig. 2 that BR($\mu \rightarrow e\gamma$) already sets the upper limit of $M_3$ below $10^{14}$ GeV.

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5 Conclusions

Abelian flavour symmetries provide interesting possibilities for understanding the hierarchy of fermion masses and mixing. Despite uncertainties in the choice of $O(1)$ coefficients, they offer useful insight into physical observables, and provide specific predictions for the signals to be expected in various detection channels, which serve as diagnostic tools for discriminating between different models; moreover, several structures predicted by non-abelian symmetries can be well reproduced by simple abelian constructions.

In our work, we have explored these possibilities, using updated experimental input from neutrino data, particularly recent measurements of $\theta_{13}$, MEG and the LHC. We have revisited the signatures of charged LFV within an SU(5) GUT framework supplemented by an abelian flavour symmetry, studying the correlations arising in CMSSM models with parameter values that are favoured by the LHC and cosmological considerations, finding interesting possibilities even within this most constrained scenario. Because of their sensitivity to flavour symmetries and model parameters that are not constrained by the neutrino data, particularly those linked to the right-handed neutrino sector, LFV searches may become a powerful tool for distinguishing between different theoretical scenarios.

We first performed a scan of different fits to the neutrino data, selecting representative fits that lead to normal neutrino hierarchies and correlations between the neutrino mixing angles that are compatible with the global analysis of neutrino data in [57]. In doing so, we paid attention to the naturalness of the fit, avoiding artificial cancellations arising from specific choices of coefficients.

We then looked at the expectations for LFV processes in the above models, identifying the range of parameters where observable signatures are possible. In general, fits with similar predictions for the neutrino parameters may lead to different LFV predictions. However, the recent input on $\theta_{13}$, combined with the new MEG bound on $\mu \rightarrow e\gamma$ as well as LHC data, does constrain the allowed structures. Further precision in the determination of neutrino parameters could lead to restrictions on the choices of model coefficients, but would not constrain the right-handed neutrino charges. New input in this respect, however, could be provided by the rates for LFV processes, since their magnitude is directly linked to these charges, unlike the neutrino mass and mixing parameters. Additional input on the right-handed neutrino sector could be obtained by requiring successful leptogenesis, which in our case can be achieved for a natural choice of parameters.

In the cases we studied, it was possible to establish correlations between the expected rates for radiative LFV decays, the LFV decay of the second lightest neutralino $\chi_2$ at the LHC and LFV in slepton decay at a future LC, for different possibilities for the structure of the heavy Majorana neutrino masses.

Within the CMSSM, the absence of a supersymmetry signal at the LHC data and the discovery of a neutral Higgs weighing $\sim 125$ GeV imply that observation of slepton flavour violation at the LHC would be difficult but possible, for points with a lighter spectrum. Observation of
LFV at the LC is also possible for the centre of mass energies above 1 TeV that are compatible with the nominal energies of CLIC. On the other hand, it should be noted here that scenarios less constrained than the CMSSM (which could fit $m_h$ with a lighter sparticle spectrum) might predict observable LFV signals at even smaller energies.

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