Predicting charged lepton flavor violation from 3-3-1 gauge symmetry

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The simplest realization of the inverse seesaw mechanism in a $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge theory offers striking flavor correlations between rare charged lepton flavor violating decays and the measured neutrino oscillations parameters. The predictions follow from the gauge structure itself without the need for any flavor symmetry. Such tight complementarity between charged lepton flavor violation and neutrino oscillations renders the scenario strictly testable.

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Preliminaries

Beyond the discovery of the Higgs boson [1, 2] no signs of genuine new physics have shown up so far at high energies. However, the existence of new physics has been established with the discovery of neutrino oscillations [3, 4], implying the existence of lepton flavor violation and nonzero neutrino masses. Unraveling the origin of the latter constitutes one of the main challenges of particle physics. While the prevailing view is that neutrino masses arise from physics associated with unification, they might as well signal novel TeV-scale physics leading to potentially large charged lepton flavor violating (LFV) rates and possibly also new phenomena testable at the LHC [5]. In this case it could well be that new physics may actually show up mainly in the form of lepton flavor violation, boosting the motivation to search for charged LFV phenomena such as the rare decay $\mu \to e\gamma$. In fact, the current limit $\text{BR}(\mu \to e\gamma) < 5.7 \times 10^{-13}$ [6] already puts severe constraints on models of new physics.

Models based on the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge theory (3-3-1) constitute a minimal extension of the Standard Model (SM) that accounts for the existence of three families of fermions, the same as the number of colors [7, 8]. They provide an economical scheme to generate tiny neutrino masses radiatively from TeV scale physics [9] and could lead to successful gauge coupling unification through neutrino masses and TeV scale physics [10]. Moreover, they naturally solve the strong $\mathcal{CP}$ problem by including in an elegant way the Peccei-Quinn symmetry [11, 12].

Here we focus on the phenomenology of lepton flavor violation in the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ schemes. For definiteness we focus on the simplest implementation of the inverse seesaw mechanism within the 3-3-1 model. We show that it offers striking flavor correlations between rare charged lepton flavor violating decays and the measured neutrino oscillations parameters. Such predictions result from the gauge theory structure itself without the need for imposing any specific flavor symmetry. We analyze the complementarity between charged LFV and neutrino oscillations, a feature which may render the 3-3-1 scenario strictly testable within the upcoming generation of LFV searches.

The Model

We consider a variant of the model introduced in [9] in which neutrinos get masses via the inverse seesaw mechanism instead of quantum corrections. The model is based on the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge symmetry, extended with a global $U(1)_\mathcal{L}$ so as to consistently define lepton number [1]. We also invoke an auxiliary parity symmetry in order to ensure a realistic quark mass spectrum. The model contains three generations of lepton triplets (ψ_L), two generations of quark triplets ($Q_L^{1,2}$), one generation of quark anti-triplet (Q_L^3), along, of course, with their iso-singlet right-handed partners, and accompanied by three generations of neutral fermion singlets (S). The gauge symmetry breaking is implemented through three scalar anti-triplets ($\phi_{1,2,3}$). The particle content of the model is summarized in table (I). The fundamental fermions interact through the exchange of 17 gauge bosons: the 8 gluons of $SU(3)_C$, the 8 "weak" W_i bosons associated to $SU(3)_L$ (4 of which form 2 electrically charged bosons, and the rest are neutral), and the B boson associated to $U(1)_X$.

The lepton representations in table (I) can be decomposed as:

$$\psi_L = \begin{pmatrix} \ell^- \\ -\nu \\ N^c \end{pmatrix}_L^{e,\mu,\tau}, \tag{1}$$

where we identify $N_L^c \equiv (\nu_R)^c$ [8].

[1] For other inverse seesaw constructions within 3-3-1 scenarios see [13, 14].

	ψ_L	l_R	$Q_L^{1,2}$	Q_L^3	U_R	t'_R	D_R	$\hat{d}_R$	S	ϕ_1	ϕ_2	ϕ_3
$SU(3)_C$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$
$SU(3)_L$	$\mathbf{3}^*$	$\mathbf{1}$	$\mathbf{3}$	$\mathbf{3}^*$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}^*$	$\mathbf{3}^*$	$\mathbf{3}^*$
$U(1)_X$	$-\frac{1}{3}$	-1	0	$+\frac{1}{3}$	$+\frac{2}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$	0	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$
$U(1)_{\mathcal{L}}$	$-\frac{1}{3}$	-1	$-\frac{2}{3}$	$+\frac{2}{3}$	0	0	0	0	$+1$	$+\frac{2}{3}$	$-\frac{4}{3}$	$+\frac{2}{3}$
$\mathbb{Z}_2$	$+$	$+$	$+$	$-$	$+$	$-$	$-$	$+$	$+$	$+$	$+$	$-$

TABLE I: Particle content of the model. Here $U_R \equiv \{u_R, c_R, t_R\}$, $D_R \equiv \{d_R, s_R, b_R\}$ and $\hat{d}_R \equiv (d'_R, s'_R)$.

In the scalar sector, on the other hand, we have:

$$\phi_1 = \begin{pmatrix} \phi_1^0 \\ -\phi_1^- \\ \tilde{\phi}_1^- \end{pmatrix}, \quad \phi_2 = \begin{pmatrix} \phi_2^+ \\ -\phi_2^0 \\ \tilde{\phi}_2^0 \end{pmatrix}, \quad \phi_3 = \begin{pmatrix} \phi_3^+ \\ -\phi_3^0 \\ \tilde{\phi}_3^0 \end{pmatrix}. \tag{2}$$

After electroweak symmetry breaking, the electric charge and lepton number assignments of the particles of the model follow from the action of the operators:

$$Q = T_3 + \frac{1}{\sqrt{3}} T_8 + X\,; \tag{3}$$

$$L = \frac{4}{\sqrt{3}} T_8 + \mathcal{L}\,. \tag{4}$$

The relevant terms in the Lagrangian for leptons are:

$$-\mathcal{L}_{\rm lep} = y^\ell \bar{\psi}_L l_R \phi_1 + y^a \overline{\psi_L^c} \psi_L \phi_1 + y^s \bar{\psi}_L S \phi_2 + \frac{m_S}{2}\, \overline{S^c} S + \text{h.c.}\,, \tag{5}$$

where y^ℓ and y^s are generic 3×3 matrices, while y^a is anti-symmetric and m_S is the 3×3 Majorana mass term for the singlets S (symmetric, due to the Pauli principle).

Scalar potential and symmetry breaking

The scalar potential of the model can be written as:

$$V = \sum_i \mu_i^2 |\phi_i|^2 + \lambda_i |\phi_i|^4 + \sum_{i \neq j} \lambda_{ij} |\phi_i|^2 |\phi_j|^2 + f\,(\phi_1 \phi_2 \phi_3 + \text{h.c.}) + m_{\rm s}^2\,(\phi_2^* \phi_3 + \text{h.c.})\,, \tag{6}$$

where $\mu_{1,2,3}$, f and $m_{\rm s}$ are parameters with dimensions of mass. The two latter couplings break the $\mathbb{Z}_2$ softly. For simplicity we denote all the dimensionless couplings by λ and take $m_{\rm s} = 0$.

In full generality, the scalars of the model are allowed to take vacuum expectation values (VEVs) in the following directions $\langle \phi_1 \rangle^T = (k_1, 0, 0)/\sqrt{2}$, $\langle \phi_2 \rangle^T = (0, k_3, n)/\sqrt{2}$, and $\langle \phi_3 \rangle^T = (0, k_2, n')/\sqrt{2}$. However, in order to recover the SM as a low energy limit, we assume the hierarchy $k_{1,2,3} \ll n, n'$. Moreover, we assume: $k_3 = n' = 0$, which together with the $\mathbb{Z}_2$ symmetry guarantees the existence of a simple pattern of realistic quark masses (see below).

We define the covariant derivative in the usual way as:

$$D_\mu = \partial_\mu - i \sum_{\rm groups} g A_\mu^a T_a\,, \tag{7}$$

where A_μ^a is the gauge boson, T_a are the generators of the group and the sum extends over all gauge groups included in $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$. Assuming $k_1 \sim k_2 \equiv k \ll n$, and keeping only the leading order terms, one finds that the mass spectrum of the charged scalars is given as: [2]

$$M^2(\phi_2^\pm) = 0\,, \tag{8}$$

$$M^2((\phi_1^\pm + \phi_3^\pm)/\sqrt{2}) = 0\,, \tag{9}$$

$$M^2((\phi_1^\pm - \phi_3^\pm)/\sqrt{2}) \sim \frac{1}{\sqrt{2}} f\,n\,, \tag{10}$$

$$M^2(\tilde{\phi}_1^\pm) \sim \sqrt{2} f\,n\,. \tag{11}$$

On the other hand the masses of the neutral CP-even scalars are, up to corrections of $\mathcal{O}(k^2)$:

$$M^2(\Re(\phi_1^0 + \phi_3^0)/\sqrt{2}) \sim (2\lambda + \sqrt{2}\,\frac{f}{n} - \frac{1}{2}\,\frac{f^2}{\lambda n^2}) k^2 \tag{12}$$

$$M^2(\Re(\phi_1^0 - \phi_3^0)/\sqrt{2}) \sim \frac{1}{\sqrt{2}} f\,n\,, \tag{13}$$

$$M^2(\Re\phi_2^0) = 0\,, \tag{14}$$

$$M^2(\Re\tilde{\phi}_2) \sim 2\lambda\,n^2\,, \tag{15}$$

$$M^2(\Re\tilde{\phi}_3) \sim \sqrt{2} f\,n\,. \tag{16}$$

Finally, the masses of the neutral CP-odd scalars at leading order are given as:

$$M^2(\Im(\phi_1^0 + \phi_3^0)/\sqrt{2}) \sim \sqrt{2} f\,n\,, \tag{17}$$

$$M^2(\Im(\phi_1^0 - \phi_3^0)/\sqrt{2}) = 0\,, \tag{18}$$

$$M^2(\Im\phi_2^0) = 0\,, \tag{19}$$

$$M^2(\Im\tilde{\phi}_2) = 0\,, \tag{20}$$

$$M^2(\Im\tilde{\phi}_3) \sim \frac{1}{\sqrt{2}} f\,n\,. \tag{21}$$

The massless scalars found in the above equations correspond to the degrees of freedom 'eaten-up' by the charged and neutral gauge bosons, respectively, which acquire the

[2] We identify the corresponding (approximate) eigenstates between parentheses.

following masses:

$$m_W^2 \;=\; \frac{1}{2}\,g_2^2\,k^2\,, \tag{22}$$

$$m_{W'}^2 \;=\; \frac{1}{4}\,g_2^2\,n^2\,, \tag{23}$$

$$m_Z^2 \;=\; \frac{g_2^2(4g_1^2+3g_2^2)}{2\,(g_1^2+3g_2^2)}\,k^2\,, \tag{24}$$

$$m_{Z'}^2 \;=\; \frac{1}{9}\,(g_1^2+3g_2^2)\,n^2\,, \tag{25}$$

$$m_X^2 \;=\; m_Y^2 \;=\; \frac{1}{4}\,g_2^2\,n^2\,. \tag{26}$$

Notice that since $\tilde{\phi}_2^0$ is singlet under the $SU(2)_L$ subgroup contained in $SU(3)_L$, the VEV n will control the four new gauge bosons masses and break $SU(3)_L$ to $SU(2)_L$. On the other hand, $SU(2)_L\otimes U(1)_Y$ is broken at the electroweak scale by the k_1 and k_2 VEVs down to the electromagnetic $U(1)_Q$ symmetry. For $f\sim n$ all the scalars of the model are naturally heavy, except one state that we can identify with the SM Higgs boson, i.e., $H\equiv(\phi_1^0+\phi_3^0)/\sqrt{2}$, in good approximation. Indeed, its couplings to the fermions confirm that the state H is the one that gives mass to SM fermions.

Quark sector

We now turn to the quark sector. From the symmetries of the model, see table (I), it follows that the quark Lagrangian is given by:

$$\begin{aligned}\mathcal{L}_{\rm quarks} \;=\; & \bar{Q}_L^{1,2}\,y^u U_R\phi_1^* + \bar{Q}_L^{1,2}\,y^d D_R\phi_3^* + \bar{Q}_L^{1,2}\,\bar{y}^d\hat{d}_R\phi_2^* \\ & + \bar{Q}_L^{3}\,\tilde{y}^u U_R\phi_3 + \bar{Q}_L^{3}\,\tilde{y}^d D_R\phi_1 + \bar{Q}_L^{3}\,\bar{y}^u t'_R\phi_2 \\ & + \text{h.c.}\,,\end{aligned} \tag{27}$$

where we defined $\hat{d}_R\equiv(d'_R,s'_R)$, $U_R\equiv\{u_R,c_R,t_R\}$ and $D_R\equiv\{d_R,s_R,b_R\}$. This Lagrangian leads to the following mass matrices:

$$M_d=-\frac{1}{\sqrt{2}}\begin{pmatrix} y^d_{11}\,k_2 & y^d_{12}\,k_2 & y^d_{13}\,k_2 & 0 & 0\\ y^d_{21}\,k_2 & y^d_{22}\,k_2 & y^d_{23}\,k_2 & 0 & 0\\ \tilde{y}^d_{11}\,k_1 & \tilde{y}^d_{12}\,k_1 & \tilde{y}^d_{13}\,k_1 & 0 & 0\\ 0 & 0 & 0 & \bar{y}^d_{14}\,n & \bar{y}^d_{15}\,n\\ 0 & 0 & 0 & \bar{y}^d_{24}\,n & \bar{y}^d_{25}\,n \end{pmatrix}, \tag{28}$$

$$M_u=-\frac{1}{\sqrt{2}}\begin{pmatrix} y^u_{11}\,k_1 & y^u_{12}\,k_1 & y^u_{13}\,k_1 & 0\\ y^u_{21}\,k_1 & y^u_{22}\,k_1 & y^u_{23}\,k_1 & 0\\ \tilde{y}^u_{11}\,k_2 & \tilde{y}^u_{12}\,k_2 & \tilde{y}^u_{13}\,k_2 & 0\\ 0 & 0 & 0 & \bar{y}^u_{14}\,n \end{pmatrix}. \tag{29}$$

Thanks to the $\mathbb{Z}_2$ symmetry, the SM and exotic subsectors are independent of each other and can be adjusted individually to easily obtain a realistic quark sector and heavy exotic quarks at the same time.

Neutrino masses and inverse seesaw mechanism

The presence of the small term $\overline{S^c}S$, in eq. (5), explicitly breaks $U(1)_{\mathcal{L}}$ and provides the seed for lepton number violation leading to neutrino masses via the inverse seesaw mechanism. Indeed, after spontaneous symmetry breaking of the electroweak gauge group, we get the following 9×9 neutrino mass matrix, in the basis (ν,N,S) [15]:

$$\mathcal{M}=\begin{pmatrix} 0 & m_D & 0\\ & 0 & M\\ & & m_S \end{pmatrix}, \tag{30}$$

where $m_D\equiv\sqrt{2}k_1\,y^a$ and $M\equiv\frac{1}{\sqrt{2}}n\,y^s$ [3]. The inverse seesaw-induced light neutrino masses can be written as [15]:

$$m_\nu = m_D\,\left(M^T\right)^{-1}\,m_S\,M^{-1}\,m_D^T\,. \tag{31}$$

Here, the matrix M can be taken diagonal without loss of generality. Using this freedom and taking into account that m_D is anti-symmetric, eq. (31) can be expressed in terms of an effective symmetric 3×3 matrix, $\widetilde{M}^{-1}\equiv M^{-1}\,m_S\,M^{-1}$, as:

$$m_\nu = -m_D\,M^{-1}\,m_S\,M^{-1}\,m_D \equiv -m_D\,\widetilde{M}^{-1}\,m_D\,. \tag{32}$$

A simple implication of the antisymmetry of the "Dirac" entry m_D is that $\mathrm{Det}(m_\nu)=0$, so that the lightest neutrino in this model must be massless at the tree level.

Lepton flavor violation predictions

Let us now proceed to a simple parameter counting. On the left-hand side of eq. (31) one has 5 independent complex parameters, since $\mathrm{Det}(m_\nu)=0$. In contrast, on the right-hand side of eq. (31) one has 9 independent complex parameters: 3 in m_D, and 6 elements in $\widetilde{M}$. Therefore, we have 4 (complex) relations among the parameters (y^s, y^a, and m_S). One can choose as free parameters the 3 off-diagonal entries of $\widetilde{M}^{-1}$, together with a global scaling factor $\tilde{m}$ defined through:

$$m_D^{ij}=\tilde{m}^{-1}\,\left(m_\nu^{1j}m_\nu^{2i}-m_\nu^{1i}m_\nu^{2j}\right)\,. \tag{33}$$

From eq. (31), we can see that $\tilde{m}$ scales as $\sqrt{m_\nu^3 m_S/M^2}$, so that for $m_S\approx 10\,\text{eV}$, $M\approx 1\,\text{TeV}$, and neutrino masses of $\mathcal{O}(0.1)\,\text{eV}$, we obtain $\tilde{m}\approx10^{-22}\,\text{GeV}$. In contrast, the

[3] Note that the matrix in eq. (30) does not depend on the conditions imposed on the VEVs. Indeed, even if $k_3\neq0$ the resulting linear seesaw term [16–18] would give only a subleading contribution $\sim m_\nu(M_W/n)^2$

diagonal entries of $\widetilde{M}^{-1}$ are functions of its off-diagonal elements and m_D. We emphasize that eq. (33) is not an *ansatz*, but the most general solution of eq. (32).

Such a parameterization makes explicit the direct relation between charged LFV observables and neutrino oscillation parameters, which is a characteristic feature of our model. Indeed, LFV in this model arises from the term:

$$-\mathcal{L}_{\rm LFV} = y^a\, \psi_L^T C^{-1} \psi_L \phi_1 + \text{h.c.}\,, \tag{34}$$

which depends solely upon the coupling y^a, hence m_D. Using eq. (33) together with $m_\nu = U_\nu^*\, m_\nu^{\rm diag}\, U_\nu^\dagger$, where U_ν is the leptonic mixing matrix in its standard parameterization in terms of three mixing angles and the Dirac phase (δ), one obtains the relevant coupling for LFV as:

$$\begin{aligned} y^a \;&=\; \begin{pmatrix} 0 & \tilde{y}^a_{12} & \tilde{y}^a_{13} \\ -\tilde{y}^a_{12} & 0 & \tilde{y}^a_{23} \\ -\tilde{y}^a_{13} & -\tilde{y}^a_{23} & 0 \end{pmatrix} \times \frac{\sqrt{\Delta m^2_{\rm atm}}}{\sqrt{2} k_1\, \tilde{m}} \\ &\times \begin{cases} \sqrt{\Delta m^2_{\rm sol}} & \text{(NH)} \\ \sqrt{\Delta m^2_{\rm atm} + \Delta m^2_{\rm sol}} & \text{(IH)} \end{cases} \,, \end{aligned} \tag{35}$$

for normal (NH) and inverse hierarchies (IH). Here the parameters $\tilde{y}^a_{ij}$ are functions of the lepton mixing matrix parameters, i.e. $\tilde{y}^a_{ij} = \tilde{y}^a_{ij}(\theta_{12}, \theta_{13}, \theta_{23}, \delta)$, and are given in the Appendix. It is remarkable that the Yukawas y^a relevant for determination of LFV rates are, up to a global scaling factor, fully determined by the parameters measured in neutrino oscillation experiments. This allows us to make definite predictions for LFV observables that can be used to provide an unambiguous test of the model, as we show below.

Radiative $\ell_i \to \ell_j \gamma$ decays

In order to show the predictive power of the model here we focus, for definiteness, on flavour-changing leptonic (radiative) decays. This probe constitutes one of the most important tests of new physics and has been actively sought after in many experiments. The branching ratio (BR) of the decay of the charged lepton $\ell_i \to \ell_j \gamma$ is given as:

$${\rm BR}(\ell_i \to \ell_j \gamma) = \frac{m^5_{\ell_i}\, |\, (y^a\, F\, y^a)_{ij}\,|^2}{\Gamma_{\ell_i}}\,, \tag{36}$$

where Γ_{ℓ_k} is the total decay width of ℓ_k, and F is a function that depends on the masses and mixings of all the particles running inside the loop (summation over the different contributions is implicit here). We have three different classes of contributions: *i*) loops mediated by the new heavy gauge bosons. These are suppressed due to the large scale of the breaking of $SU(3)_L$ compared to M_W; *ii*) contributions from the exchange of a charged scalar whose mass is $\sim \sqrt{f\,n/\sqrt{2}}$; and finally *iii*) the "standard" loop, mediated by the SM W boson and neutrinos. The latter two contributions dominate the $\ell_i \to \ell_j \gamma$ amplitude, with relative sizes depending on the ratio f/M.

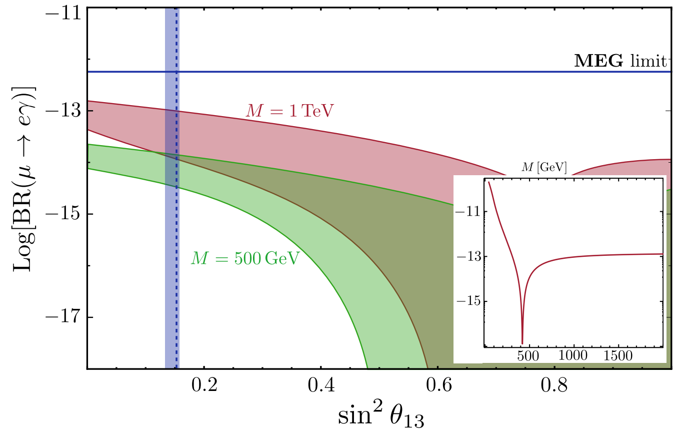

FIG. 1: The branching ratio of the decay $\mu \to e\gamma$ versus $\sin\theta_{13}$ for $M = 500\,$GeV and $1\,$TeV. We take $\tilde{m} = 2 \times 10^{-23}\,$GeV, $f = 2\,$TeV, and $n = 3\,$TeV. The vertical band is the 3σ range reported in [22]. The other mixing angles are taken within their 3σ range [22]. We also show in the lower-right corner the $\mu \to e\gamma$ branching ratio as a function of M, using best-fit values for the mixing angles, as given in [22].

This branching ratio depends on the neutrino mixing parameters, the global scaling factor $\tilde{m}$, and on the neutrino mass hierarchy. As can be seen in eq. (35), the off-diagonal entries in y^a are larger in the case of IH by a factor $\sim \sqrt{\Delta m^2_{\rm atm}/\Delta m^2_{\rm sol}}$ with respect to NH so that, for the same input parameters, one has larger LFV effects in IH.

We compute the various relevant LFV observables using `FlavorKit` [19] [4]. The branching ratio of the decay $\mu \to e\gamma$ is shown in figure (1) as a function of $\sin\theta_{13}$ for two different values of the (quasi-Dirac [8]) right-handed neutrino mass $M = 500\,$GeV and $1\,$TeV. The vertical band is the 3σ range reported in [22], whereas the other mixing angles are randomly taken within their 3σ range [22]. This figure has been obtained by varying M (by taking different values for the y^s Yukawa couplings) for the fixed parameters $n = 3\,$TeV, $f = 2\,$TeV, and $\tilde{m} = 2 \times 10^{-23}\,$GeV. We also consider degenerate right-handed neutrinos, normal hierarchy for the light

[4] This is a computer tool based on `SARAH` [20] and `SPheno` [21], that increases their capability to handle flavor observables.

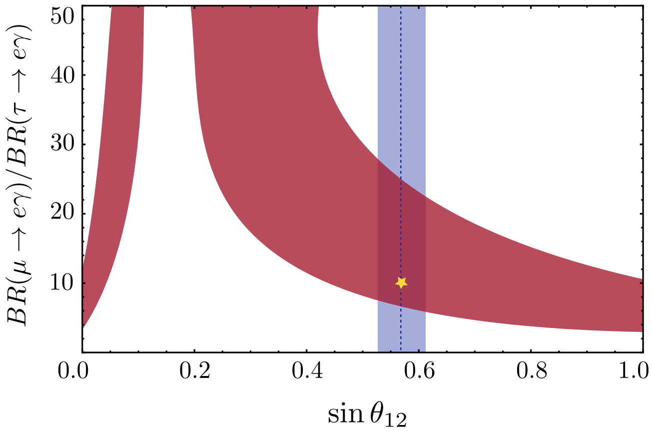

FIG. 2: Ratio of the BR of μ and τ decays to $e\gamma$, i.e., $\mathrm{BR}(\mu \to e\gamma)/\mathrm{BR}(\tau \to e\gamma)$ versus $\sin\theta_{12}$. The vertical band is the 3σ range given in [22]. The tilted band is obtained by varying the other mixing angles within their 3σ range [22].

neutrinos and a vanishing Dirac CP violating phase. One notices that the branching ratio is lower for the $M = 500\,\mathrm{GeV}$ case. This is caused by a partial cancellation between the standard loop, mediated by the W boson, and the contribution induced by the exchange of charged scalars. This cancellation takes place for $M \simeq 400\,\mathrm{GeV}$ and is explicitly illustrated in the lower-right corner of figure (1), where the $\mu \to e\gamma$ branching ratio is shown as a function of M, using best-fit values for the mixing angles [22]. The main message from figure (1) is that $\mu \to e\gamma$ may take place with sizeable rates, close to the current limit, or even larger. Given the expected sensitivities of upcoming experiments one finds that the detection of this and other muon number violating processes might become feasible.

Since the BR depends on a global multiplicative factor, it is interesting to consider the ratio of branching ratios of LFV lepton decays. It follows from eq. (36) that:

$$\frac{\mathrm{BR}(\ell_i \to \ell_j\gamma)}{\mathrm{BR}(\ell_k \to \ell_n\gamma)} = \frac{m_{\ell_i}^5}{m_{\ell_k}^5}\,\frac{|\,(y^a\,F\,y^a)_{ij}\,|^2/\Gamma_{\ell_i}}{|\,(y^a\,F\,y^a)_{kn}\,|^2/\Gamma_{\ell_k}}\,. \tag{37}$$

For the simplest case of nearly degenerate right-handed neutrinos, the F functions are all equal and cancel out in the fraction. In this case, eq. (37) depends exclusively on the ratios of $|y^a y^a|$, i.e., only on the neutrino mixing angles. The main advantage of considering the ratio of branching ratios and a quasi degenerate spectrum is that this leads to clean predictions which do not depend on the neutrino mass hierarchy nor the loop functions. Indeed, in this simplified scenario, by combining eq. (35) with eq. (37) we obtain the following predictions:

$$\frac{\mathrm{BR}(\mu \to e\gamma)}{\mathrm{BR}(\tau \to e\gamma)} = \frac{m_\mu^5\Gamma_\tau}{m_\tau^5\Gamma_\mu}\,\frac{|\tilde{y}^a_{23}\tilde{y}^a_{13}|^2}{|\tilde{y}^a_{12}\tilde{y}^a_{23}|^2} \approx 10\,, \tag{38}$$

$$\frac{\mathrm{BR}(\tau \to e\gamma)}{\mathrm{BR}(\tau \to \mu\gamma)} = \frac{|\tilde{y}^a_{12}\tilde{y}^a_{23}|^2}{|\tilde{y}^a_{12}\tilde{y}^a_{13}|^2} \approx 3\,, \tag{39}$$

where we have used the best-fit values for the neutrino parameters as derived in the global fit of neutrino oscillations given in [22] and set $\delta = 0$. So, when the right-handed neutrino spectrum is degenerate, the model predicts $\mathrm{BR}(\mu \to e\gamma) \gg \mathrm{BR}(\tau \to \ell_i\gamma)$. Therefore, given the expected sensitivities for τ LFV decays [5], the simple observation of $\tau \to \ell_i\gamma$ in one (or several) of the near future experiments would rule out our simplest degenerate right-handed neutrino hypothesis. The viable alternative scenario in such cases would be a hierarchical right-handed neutrinos spectrum, impliying a non-vanishing contribution of the F loop functions in the ratio of BRs. In this case the F functions for different flavor transitions can take very different values, and thus the ratios in Eq. (39) can clearly depart from their predictions in the degenerate scenario.

As an illustration, in figure (2) we show the ratio of the BR of τ and μ decays to $e\gamma$, namely $\mathrm{BR}(\mu \to e\gamma)/\mathrm{BR}(\tau \to e\gamma)$ as a function of the solar mixing parameter $\sin\theta_{12}$. The other oscillation parameters are varied randomly within their 3σ ranges [22]. Similarly, the ratio of the BR of leptonic τ decays, i.e., $\mathrm{BR}(\tau \to \mu\gamma)/\mathrm{BR}(\tau \to e\gamma)$ depends mainly on the solar mixing parameter.

For other LFV processes such as $\ell_i \to 3\,\ell_j$ and $\mu - e$ conversion in nuclei, our results are qualitatively similar to the ones found in standard low-scale seesaw models [24]. Loops including neutrinos give the most important contributions, leading to LFV rates comparable to the ones for the radiative decay $\ell_i \to \ell_j\gamma$. This will be of special relevance due to the expected sensitivities in the coming experiments [25]. The complete study of all LFV processes is, however, beyond the scope of this paper.

Conclusions and discussion

In summary, we have shown how a simple extended $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ electroweak gauge symmetry implementing the inverse seesaw mechanism implies striking flavor correlations between rare charged lepton flavor violating decays and the measured neutrino oscillations parameters. The predictions follow simply from the enlarged gauge structure without any imposed flavor symmetry. Such tight complementarity between charged LFV and neutrino oscillations renders the scenario strictly testable. A more detailed study of other LFV processes will be taken up elsewhere. The scheme

[5] The expected Belle II sensitivities for τ radiative decays are around 10^{-9} [23], whereas the current MEG bound on $\mathrm{BR}(\mu \to e\gamma)$ is many orders of magnitude stronger, $\mathrm{BR} < 5.7 \times 10^{-13}$.

also has a non-trivial structure in the quarks sector since, thanks to the anomaly cancelation requirements, the Glashow-Iliopoulos-Maiani mechanism breaks down, leading to a plethora of flavor-changing neutral currents in the quark sector [26, 27]. Last but not least, the model presents a rich structure of new physics at the TeV scale that could be potentially studied in the coming run of the LHC.

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Appendix: y^a Yukawa couplings

We present in this appendix the expressions for the y^a Yukawa couplings. Using the definitions of the $\tilde{y}^a$ elements in eq. (35), we find

$$\tilde{y}^a_{12} = -\left(e^{i\delta}\cos\theta_{12}\sin\theta_{13}\cos\theta_{23} - \sin\theta_{12}\sin\theta_{23}\right)^2 , \quad (40)$$

$$\tilde{y}^a_{13} = e^{i\delta}\sin\theta_{12}\cos\theta_{12}\sin\theta_{13}\cos(2\theta_{23}) - \sin\theta_{23}\cos\theta_{23}\left(\sin^2\theta_{12} - e^{2i\delta}\cos^2\theta_{12}\sin^2\theta_{13}\right) , \quad (41)$$

$$\tilde{y}^a_{23} = \cos\theta_{12}\cos\theta_{13}\left(e^{i\delta}\cos\theta_{12}\sin\theta_{13}\cos\theta_{23} - \sin\theta_{12}\sin\theta_{23}\right) . \quad (42)$$

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