Lepton flavor violation and non-unitary lepton mixing in low-scale type-I seesaw

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Within low-scale seesaw mechanisms, such as the inverse and linear seesaw, one expects (i) potentially large lepton flavor violation (LFV) and (ii) sizeable non-standard neutrino interactions (NSI). We consider the interplay between the magnitude of non-unitarity effects in the lepton mixing matrix, and the constraints that follow from LFV searches in the laboratory. We find that NSI parameters can be sizeable, up to percent level in some cases, while LFV rates, such as that for $\mu \to e\gamma$, lie within current limits, including the recent one set by the MEG collaboration. As a result the upcoming long baseline neutrino experiments offer a window of opportunity for complementary LFV and weak universality tests.

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I. INTRODUCTION

One of the biggest challenges in particle physics is to unravel the nature of the dimension five operator [1] responsible for generating the pattern of neutrino masses and mixings required to account for current oscillation experiments [2]. Within the seesaw mechanism [3, 4] this operator arises from the tree-level exchange of heavy messenger particles. For example, in the so-called type-I seesaw, these messengers are three $SU(3)\times SU(2)\times U(1)$ singlet right-handed neutrinos, which must be super-heavy in order to account for the observed smallness of neutrino masses. However, since singlets carry no gauge-anomaly, their number is theoretically unrestricted. Such generalized type-I seesaw schemes [5, 6] provide the theoretical basis for the formulation of the inverse seesaw [7] and the linear seesaw schemes [8]. Both are capable of realizing the type-I seesaw mechanism at the TeV scale and, as a result, these schemes open the possibility for novel phenomena such as

1. lepton flavor and/or CP violating processes unsuppressed by neutrino masses [9–12],
2. effectively non-unitary lepton mixing matrix [13–15] leading to non-standard effects in neutrino propagation [16–18].

Both features arise from the non-trivial structure of the electroweak currents in seesaw schemes [5, 6] [1].While their expected magnitude is negligible within the standard type-I seesaw, it can be sizeable in low-scale seesaw mechanisms.

In this paper we analyse quantitatively the interplay between these two classes of processes. More precisely we define reference parameters describing the typical magnitude of non-standard neutrino propagation effects and compare them with the constraints which arise from the searches for lepton flavor violating (LFV) processes [20–22]. For definiteness

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[1] They are generic in electroweak gauge models that mix fermions of different isospin in the weak currents [19].

we focus on two simple realizations of low-scale seesaw, namely the inverse and linear seesaw schemes. Non-unitarity effects may also arise from the charged lepton sector [23], leading also to lepton flavor violation.

We find that unitarity violation effects in neutrino propagation at the percent level are consistent with the current bounds from lepton flavor violation searches. Therefore their search at upcoming neutrino oscillation facilities [24] opens a window of opportunity to probe for new effects beyond the standard model.

The paper is organized as follows: in section II we review the type-I seesaw mechanisms, first the high-scale and then the low-scale linear and inverse seesaw realizations, fixing the notation and giving general expressions for the unitarity violating parameters; in section III we study the analytical expressions relating LFV decay branching ratios with the non-standard pieces of the general seesaw lepton mixing matrix characterizing inverse and linear seesaw schemes; in section IV we describe the parameters used in our numerical analysis, and in Sec. V we present our numerical results, both for normal and inverse neutrino mass hierarchies. Finally we give our conclusions.

II. HIGH AND LOW-SCALE SEESAW MECHANISMS

The smallness of neutrino mass with respect to that of charged fermions can be naturally accounted for within the seesaw mechanism, in which the Standard Model is extended with extra $SU(3) \times SU(2) \times U(1)$ singlets [3, 4]. In this case the resulting neutrino mass matrix in general is a $N \times N$ matrix with $N > 3$. Here we analyse the structure of this matrix in the high and low-scale type-I seesaw schemes.

A. The standard type-I seesaw

In general seesaw schemes the neutrino mass matrix M_ν can be decomposed in sub-blocks involving the standard as well as singlet neutrinos as follows

$$M_\nu = \begin{pmatrix} M_1 & M_D \\ M_D^T & M_2 \end{pmatrix}. \tag{1}$$

in the basis ν, ν^c, where the blocks $M_1 \equiv M_L$ and $M_2 \equiv M_R$ are symmetric matrices. Though the number of singlets is arbitrary, here we take an equal number of SU(2) doublets and singlets, and consider the simplest type-I seesaw, where no Higgs triplet is present, so the upper left sub-matrix $M_1 = 0$ in Eq. (1) [5] [2]. Neutrino masses arise by diagonalizing the matrix of Eq. (1),

$$U^T\, M_\nu\, U = \text{real, diagonal.} \tag{2}$$

through the transformation U connecting the weak states to the light and heavy mass eigenstates. We adopt a polar decomposition for U where

$$U = \exp(iH) \cdot V\,, \qquad H = \begin{pmatrix} 0 & S \\ S^\dagger & 0 \end{pmatrix}, \qquad V = \begin{pmatrix} V_1 & 0 \\ 0 & V_2 \end{pmatrix}. \tag{3}$$

so we have a power series expansion for Eq. (3) given as [6]:

$$U = \begin{pmatrix} \left(I - \frac{1}{2} S\, S^\dagger\right)\, V_1 & i\, S\, V_2 \\ i\, S\, V_1 & \left(I - \frac{1}{2} S^\dagger\, S\right)\, V_2 \end{pmatrix} + O(\epsilon^3) \equiv \begin{pmatrix} U_a & U_b \\ U_c & U_d \end{pmatrix}, \tag{4}$$

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[2] In this case light neutrinos get mass only as a result of the exchange of heavy gauge singlet fermions.

where we defined:

$$\epsilon \equiv M_D\, M_R^{-1}. \tag{5}$$

Substituting Eqs. (3) and (1) in Eq. (2) one finds, from the requirement of vanishing off-diagonal sub-blocks, that:

$$i\, S^* = -M_D\, M_R^{-1}. \tag{6}$$

so that, using Eq. (6) in Eq (3) we determine U as:

$$U = \begin{pmatrix} \left(I - \frac{1}{2} M_D^* (M_R^*)^{-1}\, M_R^{-1} M_D^T\right) V_1 & M_D^* (M_R^*)^{-1}\, V_2 \\ -M_R^{-1}\, M_D^T\, V_1 & \left(I - \frac{1}{2} M_R^{-1}\, M_D^T\, M_D^*\, (M_R^*)^{-1}\right) V_2 \end{pmatrix} + O(\epsilon^3), \tag{7}$$

leading to an effective light neutrino mass matrix

$$m_\nu = -M_D M_R^{-1} M_D^T\ . \tag{8}$$

This is the so called type-I seesaw mechanism. The smallness of neutrino masses follows naturally from the heaviness of the $SU(2)_L$ singlet neutrino states ν_i. Most seesaw descriptions assume equal number of doublets and singlets, $n = m = 3$. However, since singlets carry no gauge-anomaly, their number is arbitrary [5, 6]. In this paper we will consider not only the case $(n, m) = (3, 3)$ just described, but also the inverse and linear seesaw schemes, which belong to the (3,6) class, see below.

B. Inverse type-I seesaw

As an alternative to the simplest $SU(3) \times SU(2) \times U(1)$ type-I seesaw model, it has long been proposed extending the seesaw lepton content from $(3, 3)$ to $(3, 6)$, by adding three extra $SU(2)$ singlets [7] [3] S_i charged under $U(1)_L$ global lepton number the same way as the doublet neutrinos ν_i, i.e. $L = +1$. After electroweak symmetry breaking one gets the mass matrix

$$M_\nu = \begin{pmatrix} 0 & M_D & 0 \\ M_D^T & 0 & M \\ 0 & M^T & \mu \end{pmatrix}, \tag{9}$$

in the basis $\nu,\ \nu^c,\ S$, where the three ν_i^c have $L = -1$. Note that $U(1)_L$ is broken only by the nonzero $\mu_{ij} S_i S_j$ mass terms.

Generalizing the perturbative expansion method in Ref. [6] already used in the previous section one finds that the mass matrix in Eq. (9) can be block diagonalized as

$$\mathcal{U}^T \cdot M_\nu \cdot \mathcal{U} = \text{block diag} \tag{10}$$

with

$$\begin{aligned} \mathcal{U} &\approx \begin{pmatrix} I & 0 & 0 \\ 0 & \frac{1}{\sqrt{2}} I & -\frac{1}{\sqrt{2}} I \\ 0 & \frac{1}{\sqrt{2}} I & \frac{1}{\sqrt{2}} I \end{pmatrix} \begin{pmatrix} I - \frac{1}{2} S_1\, S_1^\dagger & 0 & i S_1 \\ 0 & I & 0 \\ i S_1^\dagger & 0 & I - \frac{1}{2} S_1^\dagger\, S_1 \end{pmatrix} \begin{pmatrix} I - \frac{1}{2} S_2\, S_2^\dagger & i S_2 & 0 \\ i S_2 \dagger & I - \frac{1}{2} S_2^\dagger\, S_2 & 0 \\ 0 & 0 & I \end{pmatrix}, \\ &\approx \begin{pmatrix} I & i S_2 & i S_1 \\ -i \frac{1}{\sqrt{2}} \left(S_1^\dagger - S_2^\dagger\right) & \frac{1}{\sqrt{2}} I & -\frac{1}{\sqrt{2}} I \\ i \frac{1}{\sqrt{2}} \left(S_1^\dagger + S_2^\dagger\right) & \frac{1}{\sqrt{2}} I & \frac{1}{\sqrt{2}} I \end{pmatrix} + O(\epsilon^2) \end{aligned} \tag{11}$$

[3] For simplicity we add the isosinglet pairs sequentially, though two pairs would suffice to account for the oscillations.

where S_1 and S_2 are 3×3 matrices. In the limit $\mu \to 0$ we have $S_1 = S_2 = S$ where

$$iS^* = -\frac{1}{\sqrt{2}} m_D \, (M^T)^{-1} \sim \epsilon. \tag{12}$$

With such matrix the light neutrino mass obtained after type-I seesaw is

$$m_\nu = M_D M^{T^{-1}} \mu M^{-1} M_D^T. \tag{13}$$

Note that, in the limit as $\mu \to 0$ the lepton number symmetry is recovered, making the three light neutrinos strictly massless. Thus the smallness of neutrino mass follows in a natural way, in the sense of 't Hooft [25], as it is protected by $U(1)_L$. One sees also that $\mathcal{U}$ consists of a maximal block rotation, corresponding to the Dirac nature of the three heavy leptons made-up of ν^c and S in the limit as $\mu \to 0$, and two rotations similar to Eq. (4) for the minimal type-I seesaw case considered in the previous section.

Note also that the idea behind the so-called *inverse seesaw* model can also be realized for other extended gauge groups e.g. [26–28]. Moreover, in specific models, the smallness of μ may be dynamically generated [29].

C. Linear type-I seesaw

An alternative seesaw scheme that can also be realized at low-scale is called the *linear seesaw*, and has been suggested as arising from a particular $SO(10)$ unified model [8] (for other possible constructions see [27, 30]). Once the extended gauge structure breaks down to the standard $SU(3) \times SU(2) \times U(1)$ one gets a mass matrix of the type

$$M_\nu = \begin{pmatrix} 0 & M_D & M_L \\ M_D^T & 0 & M \\ M_L^T & M^T & 0 \end{pmatrix}, \tag{14}$$

in the same basis ν, ν^c, S used in Sec. II B. Although theoretical consistency of the model requires extra ingredients, such as Higgs scalars to generate the M_L and $M \equiv M_R$ entries, here we consider just the simpler phenomenological scheme defined by the effective mass matrix in Eq. (14), as it suffices to describe the processes we are interested in.

The block-diagonalization proceeds in a very similar way as to the inverse seesaw case, in fact, for sufficiently small $M_L \, M^{-1}$ the relations Eq. (11) and Eq. (12) are the same in both schemes. One finds that the effective light neutrino mass is now given by

$$m_\nu = M_D (M_L M^{-1})^T + (M_L M^{-1}) M_D^T. \tag{15}$$

One sees that, in contrast to the "usual" seesaw relations for the effective light neutrino mass, Eqs. (8) and (13), the formula in Eq. (15) is linear in the Dirac neutrino Yukawa couplings, hence the name *linear seesaw*.

Notice also that the lepton number, defined as in the previous model, is broken only by the terms $M_L \, \nu^c \, S$. As a result one sees that, in the limit as $M_L \to 0$ the lepton number symmetry is recovered, making the three light neutrinos strictly massless. Again, as in the previous case, the smallness of neutrino mass follows in a natural way [25], as it is protected by $U(1)_L$.

III. UNITARITY VIOLATION AND THE MAGNITUDE OF LEPTON FLAVOR VIOLATION

The effective lepton mixing matrix $K_{i\alpha}$ characterizing the charged current weak interaction of mass-eigenstate neutrinos in any type of seesaw model has been fully characterized in Ref. [5]. It can be expressed in rectangular form

$$\mathcal{L} \supset i \frac{g}{\sqrt{2}} W_\nu \bar{l}_b K_{b\alpha} \gamma_\mu \nu_{\alpha L} + h.c. \, , \tag{16}$$

where

$$K_{b\alpha} = \sum_{c=1}^{n} \Omega^*_{cb} U_{c\alpha} \ , \tag{17}$$

where Ω is the 3 by 3 unitary matrix that diagonalizes the charged lepton mass matrix, while U is the unitary matrix that diagonalizes the (higher-dimensional) neutrino mass matrix characterizing the type-I seesaw mechanism of interest. We may write the K matrix as follows

$$K = (K_L, K_H) \ , \tag{18}$$

where K_L is a 3 by 3 matrix and K_H is a 3 by 6 matrix. While the rows of the K matrix are unit vectors, since $K \cdot K^\dagger = I$, the blocks K_L and K_H are not unitary.

For our purposes we can take the charged lepton mass matrix in diagonal form [4] so that $\Omega \to 1$. From Eq. (7) we obtain:

$$\begin{aligned} K_L &= \left(1 - \frac{1}{2} M_D^* (M^*)^{-1} (M)^{-1} M_D^T\right) V_1 \ , \\ K_H &= \left(M_D^* (M^*)^{-1}\right) V_2 . \end{aligned} \tag{19}$$

In order to establish a simple comparison with recent literature we parametrize the deviation from unitarity as [32]

$$K_L \equiv (1 - \eta) V_1 . \tag{20}$$

Then for the simplest high-scale type-I see-saw, the deviation from unitarity characterizing the mixing of light neutrinos, Eq. (19), is given by

$$\eta \sim \frac{1}{2} \epsilon^* \epsilon^T \approx \frac{1}{2} M_D^* (M^*)^{-1} (M)^{-1} M_D^T . \tag{21}$$

Barring *ad hoc* fine-tuning, it follows that for this case one expects negligible deviation from unitarity, namely $\epsilon \approx 10^{-10}$ and so $\eta \approx 10^{-20}$.

From now on we focus on the low-scale type-I seesaw schemes discussed in Secs II B and II C, inverse and linear seesaw, respectively. Generalizing the above discussion to these cases one finds

$$K_L = \left[I - \frac{1}{2} \left(M_D^* ((M^T)^*)^{-1} (M)^{-1} M_D^T\right)\right] V_1 \tag{22}$$

Hence the parameters characterizing the deviation from unitarity analogous to Eq. (21) are given by [5],

$$\eta^{I,L} \approx \frac{1}{2} \left(M_D^* ((M^T)^*)^{-1} (M)^{-1} M_D^T\right) , \tag{23}$$

which holds for both the type-I inverse and linear seesaw mechanisms. These parameters characterize the corresponding unitarity deviation in the light-active 3×3 sub-block of the lepton mixing matrix.

We now turn to the lepton flavor violating processes that would be induced at one loop in type-I seesaw models, as a result of the mixing of SU(2) doublet neutrinos with singlet neutral heavy leptons. The latter breaks the Glashow-Illiopoulos-Maiani cancellation mechanism [33], enhancing the rates for the loop-induced lepton flavor violating processes illustrated in Fig. 1. The $l_i \to l_j \gamma$ decay process is induced through the exchange of the nine neutral leptons coupled to the charged leptons in the charged current, namely the three light neutrinos as well as the six sub-dominantly coupled heavy states [9, 34–37]. The resulting decay branching ratio is given by

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[4] This may be automatic in the presence of suitable discrete flavor symmetries as in [31].
[5] For a recent study of unitarity violation in seesaw schemes see, for instance, Ref. [32].

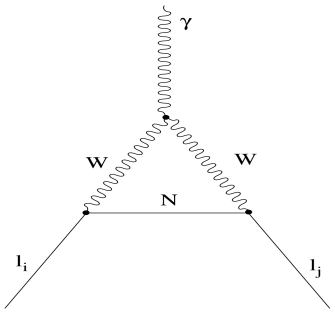

FIG. 1: Feynman graphs for $\mu \to e\gamma$ decay in seesaw models.

$$Br(l_i \to l_j \gamma) = \frac{\alpha_W^3 s_W^2}{256\pi^2} \frac{m_{l_i}^5}{M_W^4} \frac{1}{\Gamma_{l_i}} |G_{ij}^W|^2 \ , \tag{24}$$

where we use the explicit analytic form of the loop-functions [36],

$$\begin{aligned} G_{ij}^W &= \textstyle\sum_{k=1}^{9} K_{ik}^* K_{jk} G_\gamma^W \left(\frac{m_{N_k}^2}{M_W^2} \right) \ , \\ G_\gamma^W(x) &= \tfrac{1}{12(1-x)^4} (10 - 43x + 78x^2 - 49x^3 + 18x^3 \ln x + 4x^4) \end{aligned} \tag{25}$$

and is presented in Fig. 2 for each case. The difference between the two models follows from the different dependence with the lepton number violating parameters that characterize these two low-energy type-I seesaw realizations.

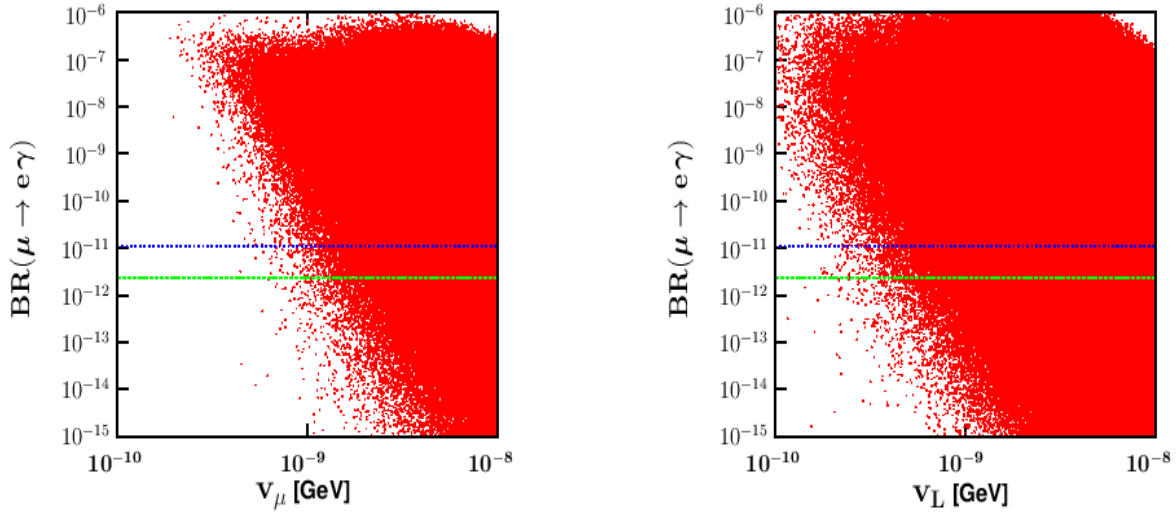

FIG. 2: $Br(\mu \to e\gamma)$ versus the lepton number violation scale: v_μ for the inverse seesaw (left panel), and v_L for the linear seesaw (right panel). In both cases one assumes normal hierarchy and the parameters are varied as explained in Sec. IV. We also indicate the limits from the MEGA collaboration [22] (upper line) and the recent limit from MEG [21] (lower line).

One sees that the branching rations may easily exceed current limits. This reflects an important feature of low-scale seesaw models, namely, that lepton flavor violation as well as leptonic CP violation proceed even in the limit of massless neutrinos [9–12]. Unsuppressed by the smallness of neutrino mass, the expected rates are sizeable. The corresponding radiative LFV decays $\tau \to e\gamma$ and $\tau \to \mu\gamma$ are not as constraining as $\mu \to e\gamma$. Similarly, the fact that the neutral current couplings of charged leptons is flavor diagonal implies that the decay processes $l_i \to l_j l_k l_r$ are suppressed by a factor of α_{QED} with respect to the radiative processes, hence less restrictive.

Another very important lepton flavor violating process is mu-e conversion in nuclei, which arises both from short range (non-photonic) as well as long range (photonic) contributions. Explicit calculations [38] indicate that, given the nuclear form factors, one finds that current mu-e conversion sensitivities are effectively lower than those of the $\mu \to e\gamma$ decay. However the upcoming generation of nuclear conversion experiments aims at substantial improvement.

IV. NUMERICAL ANALYSIS

In order to perform our numerical calculations it is convenient to generalize the Casas-Ibarra parametrization [39] to the inverse and linear type-I seesaw schemes. For simplicity we will assume real lepton Yukawa couplings and mass entries.

A. Inverse type-I Seesaw

First note that one has always the freedom to go to the basis where the 3×3 gauge-singlet block M is taken diagonal. For real m_D matrix elements, we have in total 18 parameters, nine characterizing m_D, three characterizing M, plus six from the μ matrix.

The Dirac neutrino mass matrix may be rewritten as

$$m_D = V_1 \, \mathrm{diag}(\sqrt{\tilde{m}_i}) \, R^T \, (\sqrt{\mu})^{-1} \, \mathrm{diag}(M_i^T) \tag{26}$$

where V_1 is (approximately) the mixing matrix determined in oscillation experiments [2], $\tilde{m}_i$ are the three light neutrino masses. On the other hand the arbitrary real orthogonal 3×3 matrix R and the arbitrary 3×3 real matrix M are parameters characterizing the model. This parametrization for the inverse type-I seesaw is similar to that given in Ref. [35].

In order to further reduce the number of degrees of freedom, we will make the "minimal flavor violation hypothesis" [6] which consists in assuming that flavor is violated only in the "standard" Dirac Yukawa coupling. Under this simplification the 3×3 matrix μ must be also diagonal, reducing the parameter count from a total of 18 down to 15. These include the three light neutrino masses, and the three neutrino mixing angles contained in V_1. Next come the nine model-defining parameters, that may be taken as three parameters from the R matrix, three from the μ matrix, plus three parameters characterizing M.

We have performed a scan at 3σ over the lightest mass and oscillation parameters, the three angles and the three masses in V_1 and $\tilde{m}_i$, respectively. For the scan over oscillation parameters we have used the 3σ determinations given in [2], and for the lightest mass parameter we took the cosmological bound from [43]. We parametrize the real orthogonal matrix R as a product of three rotations, marginalizing over the three angles from $0 - 2\pi$ values.

We have also fixed the upper value of the Dirac mass matrix to $(m_D)_{ij} < 175\, GeV$ to be consistent with perturbativity of the theory. The remaining six free parameters, from $diag\{\mu_{ii}\}$ and $diag\{M_{ii}\}$ matrices, are scanned as a perturbation from the identity matrix in the following way:

$$\begin{aligned} \mu_{ii} &= v_\mu \, (1 + \varepsilon_{ii}) \\ M_{ii} &= v_M \, (1 + \varepsilon'_{ii}) \, , \end{aligned} \tag{27}$$

where $|\varepsilon| \sim 5 \times 10^{-1}$. The parameter v_M setting the M-scale was fixed to $1\, TeV$, while v_μ scale was scanned in the range $(0.1 - 10)\, eV$. The two scales $v_{\mu,M}$ are consistent with the observed neutrino masses.

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[6] This simplifying assumption suffices to illustrate the points made in this paper. For alternative minimal flavor violation definitions see Refs. [40–42].

B. Linear type-I Seesaw

Similarly, for the linear seesaw case we can parametrize Dirac neutrino mass matrix as follows,

$$m_D = V_1 \text{diag}\{\sqrt{m_i}\}\, A^T \text{diag}\{\sqrt{m_i}\}\, V_1^T \left(M_L^T\right)^{-1}\, M^T, \tag{28}$$

where A has the following general form:

$$\begin{pmatrix} \frac{1}{2} & a & b \\ -a & \frac{1}{2} & c \\ -b & -c & \frac{1}{2} \end{pmatrix}, \tag{29}$$

with a, b, c real numbers. In this case M_L and $M \equiv M_R$ are general real matrices. One can always go to a basis where one of them is diagonal, for example M_L, reducing the total number of model parameters to 21.

In order to further reduce the number of degrees of freedom, we make a similar "minimal flavor violation hypothesis" to this scheme too, namely, we choose the M matrix also to be diagonal, reducing from 21 to 15 parameters. The most general Dirac mass matrix is parametrized as in Eq. (28). Hence we have a model with 15 parameters that may be taken as three parameters from A matrix, three from the M_L matrix, three from the matrix M, in addition to the three light neutrino masses and three mixing parameters.

As for the inverse type-I seesaw case, we could do the scan over the light neutrino mass and oscillation parameters and the 9 free parameters. The difference in this case respect to the inverse, is the structure of the A matrix in Eq. (29). We scan over the A matrix parameters in the form:

$$A_{ij}\ \epsilon\, (0 - 10^{-2}), \tag{30}$$

and now we define M_{Lii} in analogy to M_{ii} in Eq. (27) and we vary $v_L = (0.01 - 10)\, eV$.

V. NUMERICAL RESULTS

Low-scale seesaw schemes lead to sizeable rates for lepton flavor violating processes as well as to non-standard effects in neutrino propagation associated to non-unitary lepton mixing matrix. In this section we will quantify the interplay between these, more precisely, between the branching ratio of Eq. (24) in the low-scale type-I seesaw schemes considered here and the magnitude of the unitarity deviation defined in Eq. (21), taking into account Eqs. (26) and (28) and the requirement of acceptable light neutrino masses. For definiteness we assume leptonic CP conservation so that all lepton Yukawa couplings and mass entries are real.

We have computed the branching ratio (BR) for the charged lepton flavor violating radiative processes using Eq. (24), accurate to order $O(\epsilon^3)$ in the neutrino diagonalizing matrix, and displayed the degree of correlation of these observables with the corresponding unitarity violating parameters $|\eta_{ij}|$.

In Fig. 3 and Fig. 4 we show the results for the inverse type-I seesaw scheme with Normal (NH) and Inverted (IH) neutrino mass hierarchy for the process $\mu \to e\gamma$, respectively. The points result from a scan over the light neutrino mass and mixing parameters at 3 σ, together with the scan over the 9 free model parameters defined above.

In Fig. 5 and Fig. 6 we show the corresponding results for the $\tau \to e\gamma$ branching ratio in the inverse type-I seesaw scheme with Normal Hierarchy (NH) and Inverted Hierarchy (IH), respectively.

In Fig. 7 and Fig. 8 we show the corresponding results for the process $\tau \to \mu\gamma$ within the inverse type-I seesaw, for Normal Hierarchy (NH) and Inverted Hierarchy (IH) neutrino spectra, respectively.

Now we turn to the linear type-I seesaw scheme. In Fig. 9 and Fig. 10 we show our results for the branching ratio for the process $\mu \to e\gamma$ in such linear seesaw for Normal Hierarchy (NH) and Inverted Hierarchy (IH), respectively. The points are obtained through a scan over the neutrino oscillation parameters, as well as the free model parameters, as already described.

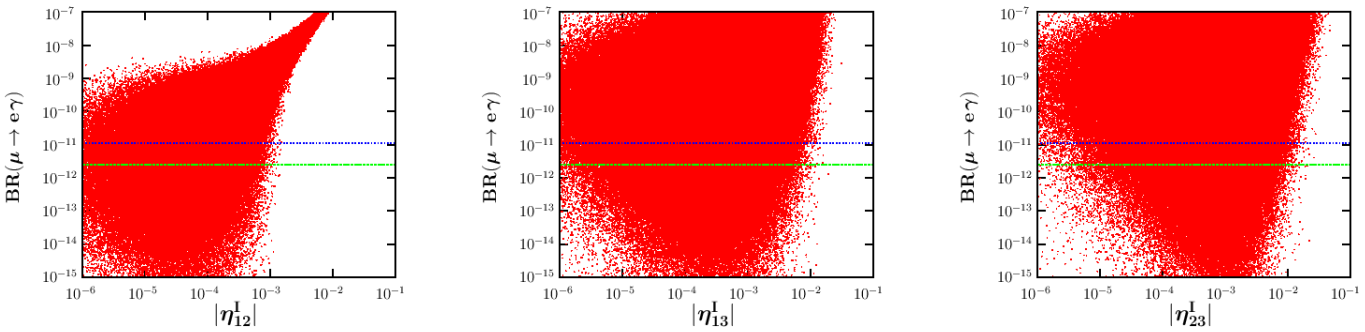

FIG. 3: Branching for the process $\mu \to e\gamma$ in type-I inverse seesaw scheme with normal hierarchy (NH). We scan over the light neutrino mass and mixing parameters at 3 σ, and over the model parameters, fixing v_M at 1 TeV and varying the v_μ scale from $1 \times 10^{-10}\, GeV$ to $1 \times 10^{-8}\, GeV$. We also indicate the limits from the MEGA collaboration (upper line) and the recent limit from MEG (lower line).

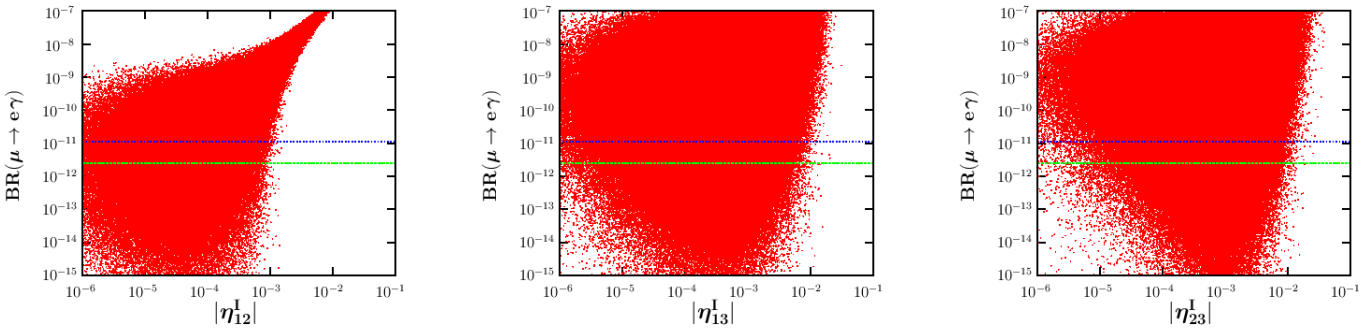

FIG. 4: Same as Fig. 3 for the Inverted Hierarchy (IH) case.

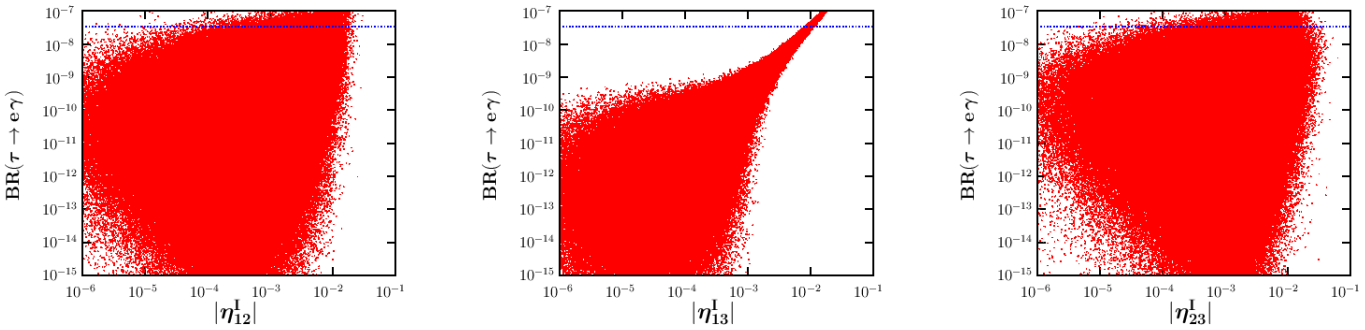

FIG. 5: $\tau \to e\gamma$ branching ratio in the inverse type-I seesaw with Normal Hierarchy (NH). The scan is performed as in Fig. 3 and the indicated limit is from the PDG [20].

In Fig. 11 and Fig. 12 we show our results for the $\tau \to e\gamma$ branching ratio, for the Normal Hierarchy (NH) and Inverted Hierarchy (IH) case, respectively. Finally, in Fig. 13 and Fig. 14 we present the corresponding results for the $\tau \to \mu\gamma$ process.

These results are summarized in table I. One sees that the magnitude of non-unitarity effects in the lepton mixing matrix can reach up to percent level is not in conflict with the constraints that follow from lepton flavor violation searches in the laboratory. Given the large - TeV scale - assumed masses of the singlet "right-handed" neutrinos there are no direct search constraints [44–46]. The main factor limiting the magnitude of non-unitarity effects then becomes the weak universality constraints. As expected, there is stronger degree of correlation between $\mu \to e\gamma$ and η_{12} than

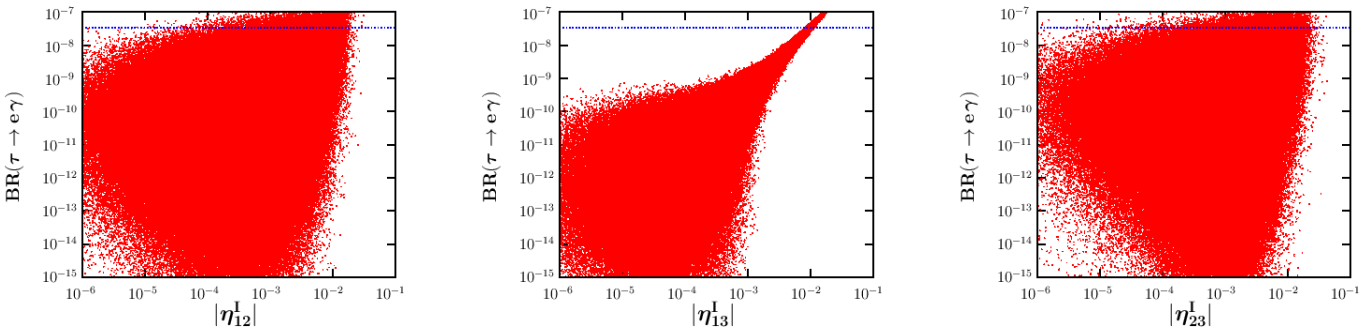

FIG. 6: Same as Fig. 5 for Inverted Hierarchy (IH).

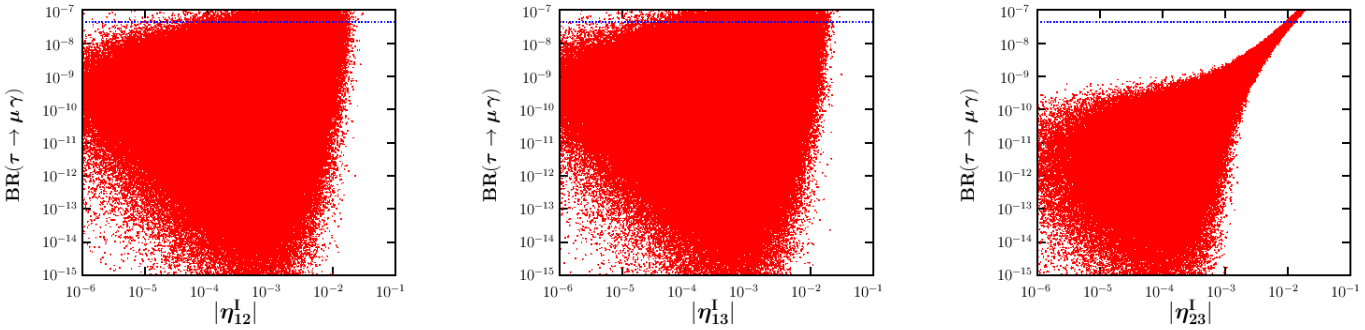

FIG. 7: $\tau \to \mu\gamma$ branching ratio in the inverse type-I seesaw with Normal Hierarchy (NH). The scan is performed as in Fig. 3 and the indicated limit is from the PDG [20].

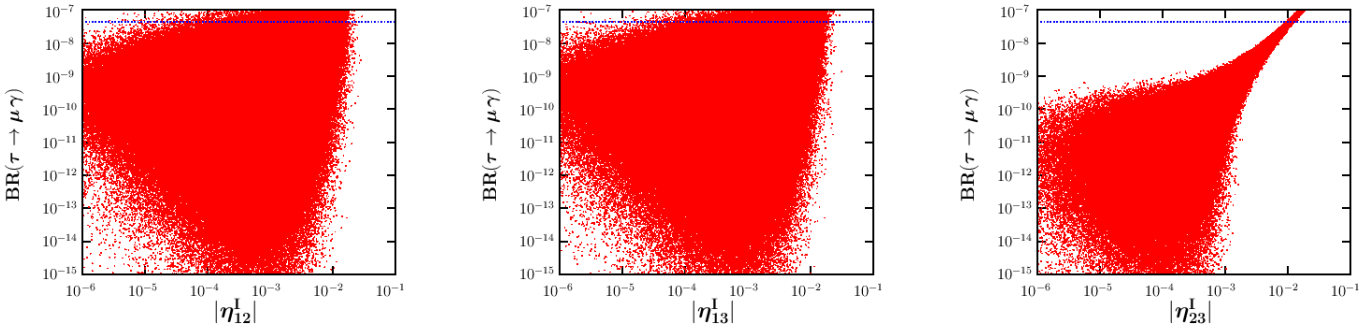

FIG. 8: Same as Fig. 7 for Inverted Hierarchy (IH).

other η's, or between $\tau \to e\gamma$ and η_{13} than others, etc. As a result in these cases one obtains the strongest restriction on unitarity violation.

Before closing let us comment on the robustness or our results with respect to the assumptions made in Sec. IV. Insofar as the regions obtained in Figs. (3)-(14) are concerned, we can state that they remain "allowed" once one departs from our simplifying assumptions. As expected however, we have verified that the regions obtained away from the simplifying assumptions may allow for somewhat larger values of the lepton-flavor-violating parameters η affecting neutrino propagation. However, on account of weak universality constraints, we prefer to stick to the more conservative values we have presented in Table I.

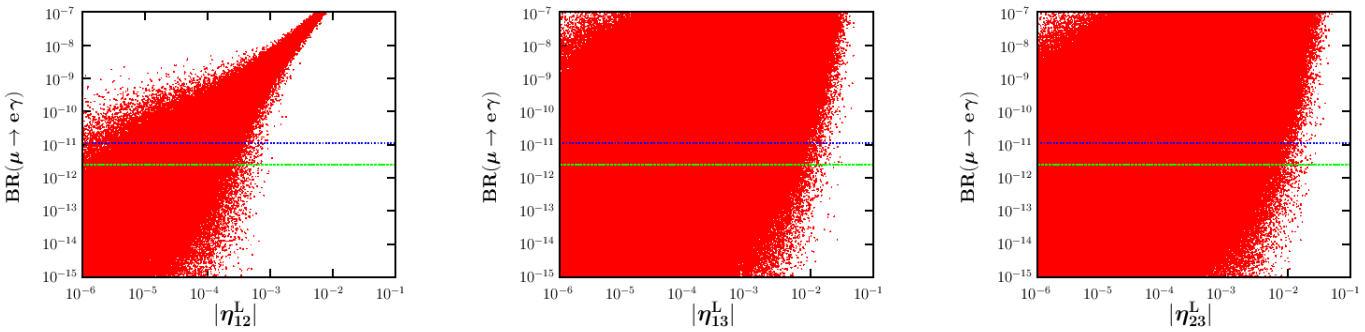

FIG. 9: Branching for the process $\mu \to e\gamma$ in type-I linear seesaw with Normal Hierarchy (NH). We scan over the parameters as in Fig. 3. The limits from the MEGA and MEG collaborations are indicated by the upper and lower horizontal lines.

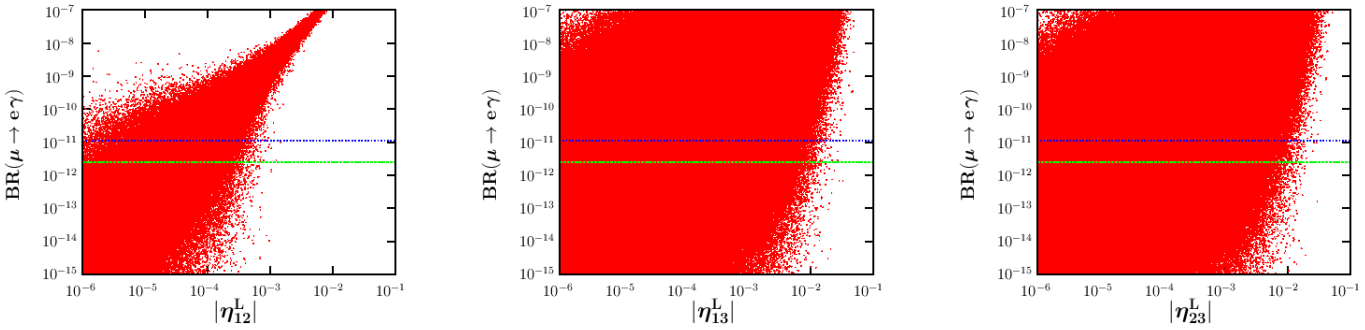

FIG. 10: Same as in Fig. 9 for Inverted Hierarchy (IH).

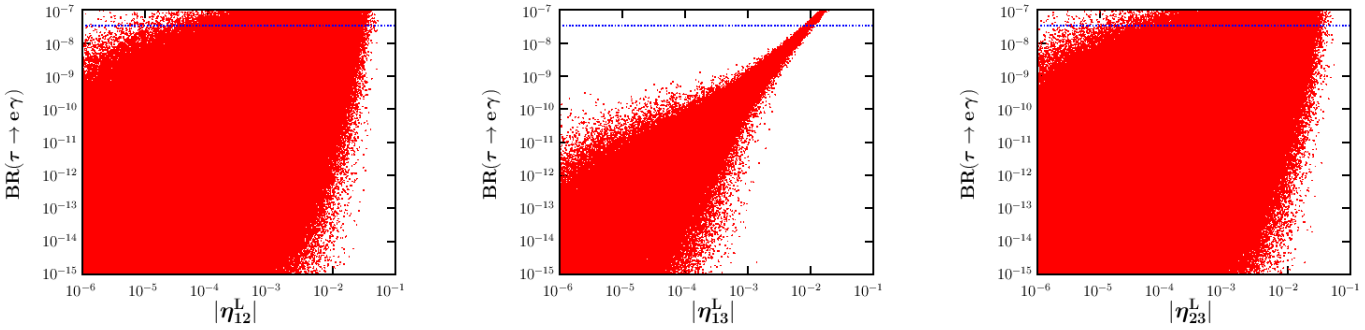

FIG. 11: $\tau \to e\gamma$ branching ratio, in Linear type-I seesaw for Normal Hierarchy (NH). Parameters scanned as in Fig. 3 and the indicated limit is from the PDG [20].

VI. CONCLUSIONS

The physics responsible for neutrino masses could lie at the TeV scale. In this case it is very unlikely that neutrino masses are not accompanied by non-standard neutrino interactions that could reveal novel features in neutrino propagation. Similarly, lepton flavor violation should also take place in processes involving the charged leptons.

Within low-scale seesaw mechanisms, such as the inverse and linear type-I seesaw, we have found that non-unitarity in the lepton mixing matrix up to the percent level in some cases, is consistent with the constraints that follow from lepton flavor violation searches in the laboratory. This conclusion holds even within the simple "minimal flavor violation" assumptions we have made. As a result the upcoming long baseline neutrino experiments [24] do provide an important window of opportunity to perform complementary tests of lepton flavor violation in neutrino propagation

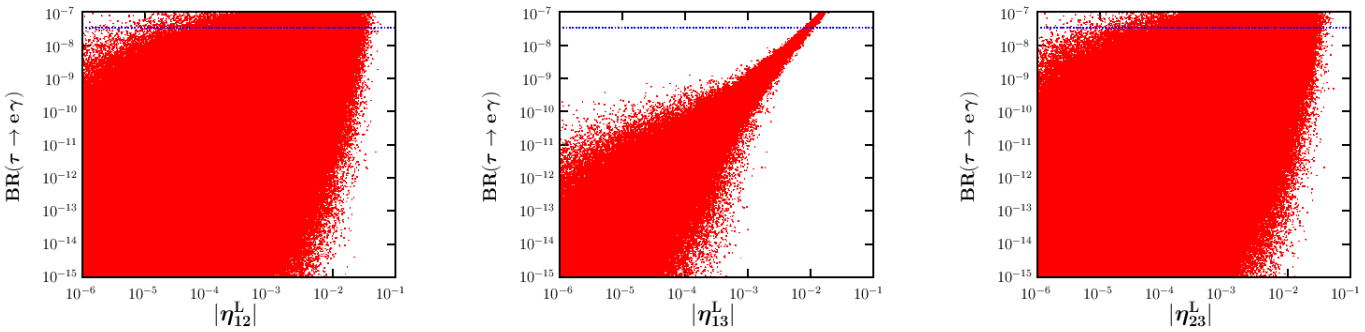

FIG. 12: Same as Fig. 11 for Inverted Hierarchy (IH).

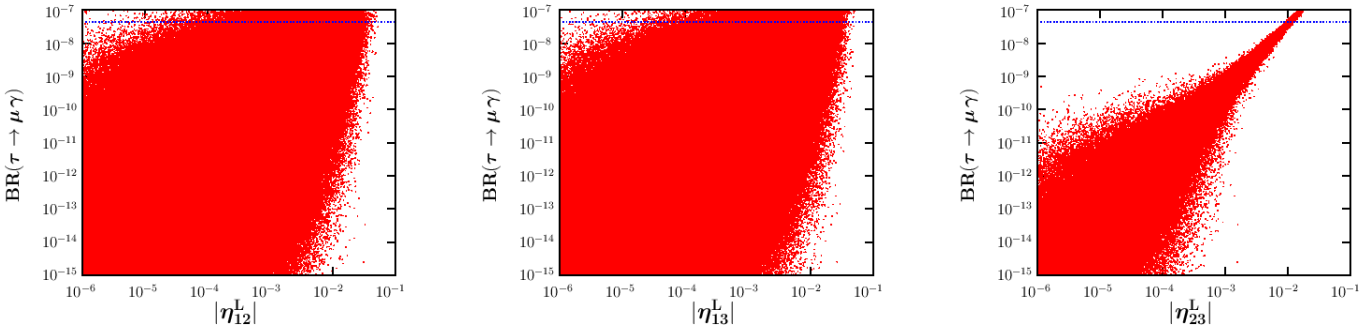

FIG. 13: $\tau \to \mu\gamma$ branching ratio, in Linear type-I seesaw for Normal Hierarchy (NH). Parameters scanned as in Fig. 3 and the indicated limit is from the PDG [20].

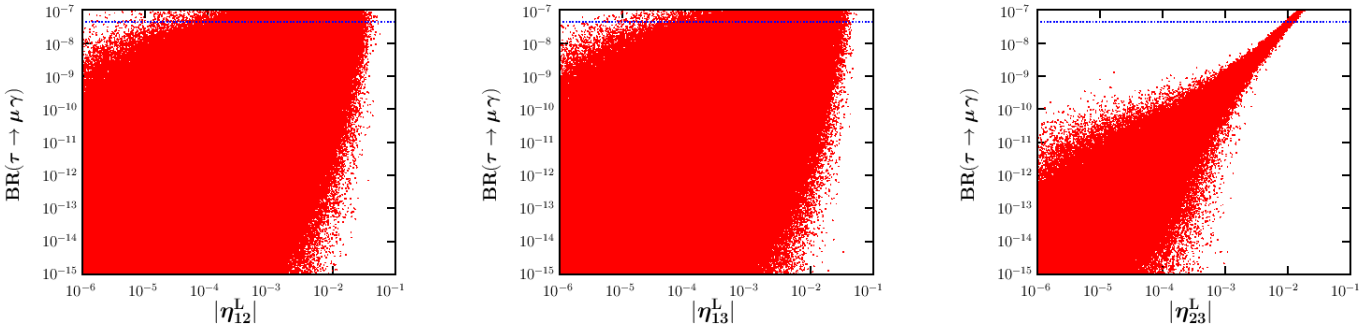

FIG. 14: Same as Fig. 11 for Inverted Hierarchy (IH).

and to probe the mass scale characterizing the seesaw mechanism.

Process	$\mu \to e\gamma$		$\tau \to e\gamma$		$\tau \to \mu\gamma$	
Hierarchy	NH	IH	NH	IH	NH	IH
$\lvert\eta^I_{12}\rvert <$	1.4×10^{-3}	1.4×10^{-3}	2.8×10^{-2}	2.8×10^{-2}	2.8×10^{-2}	2.8×10^{-2}
$\lvert\eta^I_{13}\rvert <$	2.0×10^{-2}	$2.1(1.6) \times 10^{-2}$	1.1×10^{-2}	1.1×10^{-2}	3.1×10^{-2}	3.2×10^{-2}
$\lvert\eta^I_{23}\rvert <$	$2.7(2.1) \times 10^{-2}$	$2.5(1.9) \times 10^{-2}$	6.4×10^{-2}	4.3×10^{-2}	1.2×10^{-2}	1.2×10^{-2}
$\lvert\eta^L_{12}\rvert <$	$11.0(9.6) \times 10^{-4}$	$1.5(1.1) \times 10^{-3}$	5.1×10^{-2}	5.2×10^{-2}	5.3×10^{-2}	5.7×10^{-2}
$\lvert\eta^L_{13}\rvert <$	$3.1(2.7) \times 10^{-2}$	3.3×10^{-2}	1.1×10^{-2}	1.0×10^{-2}	4.8×10^{-2}	4.8×10^{-2}
$\lvert\eta^L_{23}\rvert <$	$2.8(2.2) \times 10^{-2}$	3.0×10^{-2}	5.5×10^{-2}	5.4×10^{-2}	1.2×10^{-2}	1.2×10^{-2}

TABLE I: Limits on unitarity violation parameters from lepton flavor violation searches. The numbers given in parenthesis correspond to the improvement obtained with the recent MEG limit on $\mu \to e\gamma$. Other entries in the table are unchanged. These limits express the correlation between lepton non-unitarity and LFV that holds in low-scale seesaw schemes under a "minimal flavor violation hypothesis" defined in the text.

VII. ACKNOWLEDGMENTS

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[1] S. Weinberg, Phys. Rev. **D22**, 1694 (1980).
[2] T. Schwetz, M. Tortola, and J. W. F. Valle, New J. Phys. **13**, 063004 (2011), latest T2K/MINOS results in addendum in arXiv:1108.1376 [hep-ph] (to be published).
[3] R. Mohapatra, Talk at Neutrino 2010 Conference, Athens.
[4] J. W. F. Valle, J. Phys. Conf. Ser. **53**, 473 (2006), Review lectures at Corfu, hep-ph/0608101.
[5] J. Schechter and J. W. F. Valle, Phys. Rev. **D22**, 2227 (1980).
[6] J. Schechter and J. W. F. Valle, Phys. Rev. **D25**, 774 (1982).
[7] R. N. Mohapatra and J. W. F. Valle, Phys. Rev. **D34**, 1642 (1986);
M. C. Gonzalez-Garcia and J. W. F. Valle, **B216**, 360 (1989).
[8] M. Malinsky, J. C. Romao, and J. W. F. Valle, Phys. Rev. Lett. **95**, 161801 (2005).
[9] J. Bernabeu *et al.*, Phys. Lett. **B187**, 303 (1987).
[10] M. C. Gonzalez-Garcia and J. W. F. Valle, Mod. Phys. Lett. **A7**, 477 (1992).
[11] G. C. Branco, M. N. Rebelo, and J. W. F. Valle, Phys. Lett. **B225**, 385 (1989).
[12] N. Rius and J. W. F. Valle, Phys. Lett. **B246**, 249 (1990).
[13] S. Antusch, J. P. Baumann, and E. Fernandez-Martinez, Nucl. Phys. **B810**, 369 (2009), 0807.1003.
[14] M. Malinsky, T. Ohlsson, and H. Zhang, Phys. Rev. **D79**, 073009 (2009), 0903.1961.
[15] P. Dev and R. Mohapatra, Phys.Rev. **D81**, 013001 (2010), 0910.3924.
[16] J. W. F. Valle, Phys. Lett. **B199**, 432 (1987).
[17] H. Nunokawa *et al.*, Phys. Rev. **D54**, 4356 (1996);
A. Esteban-Pretel, R. Tomas, and J. W. F. Valle, Phys. Rev. **D76**, 053001 (2007).
[18] O. G. Miranda, M. A. Tortola, J. W. F. Valle, JHEP **0610**, 008 (2006);
N. Fornengo, M. Maltoni, R. Tomas, J. W. F. Valle, Phys. Rev. **D65**, 013010 (2002).
[19] B. W. Lee and R. E. Shrock, Phys. Rev. **D16**, 1444 (1977).
[20] K. Nakamura *et al.*, Journal of Physics G: Nuclear and Particle Physics **37**, 075021 (2010).
[21] MEG collaboration, J. Adam *et al.*, (2011), 1107.5547.
[22] MEGA collaboration, M. Ahmed *et al.*, Phys. Rev. **D65**, 112002 (2002), hep-ex/0111030.

[23] D. Ibanez, S. Morisi, and J. W. F. Valle, Phys. Rev. **D80**, 053015 (2009), [arXiv:0907.3109 [hep-ph]]; J. F. Kamenik, M. Nemevsek, JHEP **0911**, 023 (2009), [arXiv:0908.3451 [hep-ph]].

[24] ISS Physics Working Group, A. Bandyopadhyay *et al.*, Rep. Prog. Phys. **72**, 106201 (2009).

[25] G. t'Hooft, Lectures at Cargese Summer Inst. 1979, p.345-367, Ed. by E. Farhi et al. (World Scientific, Singapore)

[26] D. Wyler and L. Wolfenstein, Nucl. Phys. **B218**, 205 (1983).

[27] E. Akhmedov *et al.*, Phys. Rev. **D53**, 2752 (1996), hep-ph/9509255.

[28] S. M. Barr and I. Dorsner, Phys. Lett. **B632**, 527 (2006), hep-ph/0507067.

[29] F. Bazzocchi et. al., Phys.Rev. **D81**, 051701 (2010), 0907.1262.

[30] E. Akhmedov *et al.*, Phys. Lett. **B368**, 270 (1996), hep-ph/9507275.

[31] M. Hirsch, S. Morisi, and J. Valle, Phys. Lett. B **679**, 454 (2009).

[32] H. Hettmansperger, M. Lindner, and W. Rodejohann, arXiv:1102.3432v1 [hep-ph].

[33] S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. **D2**, 1285 (1970).

[34] A. Ilakovac and A. Pilaftsis, Nucl. Phys. **B437**, 491 (1995), hep-ph/9403398.

[35] F. Deppisch and J. W. F. Valle, Phys. Rev. **D72**, 036001 (2005), hep-ph/0406040.

[36] B. He, T. Cheng, and L.-F. Li, Phys.Lett. **B553**, 277 (2003), hep-ph/0209175.

[37] L. Lavoura, Eur. Phys. J. **C29**, 191 (2003), hep-ph/0302221.

[38] F. Deppisch, T. S. Kosmas, and J. W. F. Valle, Nucl. Phys. **B752**, 80 (2006), hep-ph/0512360.

[39] J. A. Casas and A. Ibarra, Nucl. Phys. **B618**, 171 (2001), hep-ph/0103065.

[40] G. D'Ambrosio, G. Giudice, G. Isidori, and A. Strumia, Nucl.Phys. **B645**, 155 (2002), hep-ph/0207036.

[41] V. Cirigliano, B. Grinstein, G. Isidori, and M. B. Wise, Nucl.Phys. **B728**, 121 (2005), hep-ph/0507001.

[42] R. Alonso, G. Isidori, L. Merlo, L. A. Munoz, and E. Nardi, JHEP **1106**, 037 (2011), 1103.5461.

[43] G. L. Fogli, E. Lisi, A. Marrone, A. Melchiorri, A. Palazzo, A. M. Rotunno, P. Serra, J. Silk, and A. Slosar, Phys. Rev. D **78**, 033010 (2008).

[44] M. Dittmar, A. Santamaria, M. C. Gonzalez-Garcia, and J. W. F. Valle, Nucl. Phys. **B332**, 1 (1990).

[45] S. Kovalenko, Z. Lu, and I. Schmidt, Phys.Rev. **D80**, 073014 (2009), 0907.2533.

[46] A. Atre, T. Han, S. Pascoli, and B. Zhang, JHEP **0905**, 030 (2009), 0901.3589.