

Higgs decay $h \rightarrow \mu\tau$ with minimal flavor violation

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ABSTRACT: We consider the tentative indication of flavor-violating Higgs boson decay $h \rightarrow \mu\tau$ recently reported in the CMS experiment within the framework of minimal flavor violation. Specifically, we adopt the standard model extended with the seesaw mechanism involving right-handed neutrinos plus effective dimension-six operators satisfying the minimal flavor violation principle in the lepton sector. We find that it is possible to accommodate the CMS $h \rightarrow \mu\tau$ signal interpretation provided that the right-handed neutrinos couple to the Higgs boson in some nontrivial way. We take into account empirical constraints from other lepton-flavor-violating processes and discuss how future searches for the $\mu \rightarrow e\gamma$ decay and $\mu \rightarrow e$ conversion in nuclei may further probe the lepton-flavor-violating Higgs couplings.

KEYWORDS: Higgs Physics, Beyond Standard Model, Neutrino Physics, Global Symmetries

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1 Introduction

The Higgs boson discovered at the LHC three years ago [1, 2] can offer a potential window into physics beyond the standard model (SM). The existence of new interactions can bring about modifications to the standard decay modes of the particle and/or cause it to undergo exotic decays [3]. As LHC data continues to accumulate with increasing precision, they may reveal clues of new physics in the Higgs couplings.

The latest LHC measurements of the Higgs, h , have started to expose its Yukawa interactions with leptons. Particularly, the ATLAS and CMS Collaborations have observed the decay mode $h \rightarrow \tau^+\tau^-$ and measured its signal strength to be $\sigma/\sigma_{\text{SM}} = 1.44^{+0.42}_{-0.37}$ and 0.91 ± 0.28 , respectively [4, 5]. In contrast, their direct searches for the decay channel $h \rightarrow \mu^-\mu^+$ have so far come up with only upper limits on its branching fraction, $\mathcal{B}(h \rightarrow \mu^-\mu^+) < 1.5 \times 10^{-3}$ and 1.6×10^{-3} , respectively [6, 7], at 95% confidence level (CL). Overall, these results are still consistent with SM expectations.

There have also been searches for flavor-violating dilepton Higgs decays, which the SM does not accommodate. In this regard, CMS recently reported [8] the interesting detection of a slight excess of $h \rightarrow \mu^\pm\tau^\mp$ events with a significance of 2.4σ . If interpreted as a signal, the excess implies a branching fraction of $\mathcal{B}(h \rightarrow \mu\tau) = \mathcal{B}(h \rightarrow \mu^-\tau^+) + \mathcal{B}(h \rightarrow \mu^+\tau^-) = (0.84^{+0.39}_{-0.37})\%$, but as a statistical fluctuation it translates into the bound $\mathcal{B}(h \rightarrow \mu\tau) < 1.51\%$ at 95% CL [8]. In view of its low statistical significance, it is too soon to draw a definite conclusion from this finding, but it would constitute evidence of new physics if confirmed by future experiments.

This tantalizing, albeit tentative, hint of lepton flavor violation (LFV) outside the neutrino sector has attracted a growing amount of attention, as the detection of such a process would serve as a test for many models [9–22] and could have major implications for upcoming Higgs measurements [22–28]. Subsequent to the $h \rightarrow \mu\tau$ announcement by CMS, its signal hypothesis was theoretically examined in the contexts of various scenarios involving enlarged scalar sectors [29–43] or nonrenormalizable effective interactions [39–44].

In this paper, we follow the latter line of approach which relies on effective operators to address LFV in Higgs decay. To handle the LFV pattern systematically without getting into model details, we adopt the framework of so-called minimal flavor violation (MFV). Motivated by the fact that the SM has succeeded in describing the existing data on flavor-changing neutral currents and CP violation in the quark sector, the MFV principle presupposes that Yukawa couplings are the only sources for the breaking of flavor and CP symmetries [45–50]. However, unlike its straightforward implementation for quarks, there is no unique way to extend the notion of MFV to leptons, as the minimal version of the SM by itself, without right-handed neutrinos or extra scalar particles, does not accommodate LFV. In light of the fact that flavor mixing among neutrinos has been empirically established [51], it is attractive to formulate leptonic MFV by incorporating new ingredients that can explain this observation [52]. Thus, here we consider the SM expanded with the addition of three heavy right-handed neutrinos as well as effective dimension-six operators conforming to the MFV criterion.¹ The heavy neutrinos are essential for the seesaw mechanism to endow light neutrinos with Majorana masses.

In the next section, after briefly reviewing the MFV framework, we introduce the effective dimension-six operators that can give rise to LFV in Higgs decay, only one of which is relevant to $h \rightarrow \mu\tau$. In section 3, we explore the parameter space associated with this operator which can yield $\mathcal{B}(h \rightarrow \mu\tau) \sim 1\%$, as CMS may have discovered. At the same time, we take into account various experimental restrictions on the Higgs couplings proceeding from the operator. Specifically, we impose constraints inferred from the LHC measurements described above as well as from the existing data on transitions with LFV that have long been the subject of intensive quests, such as $\mu \rightarrow e\gamma$. We present several sample points from the viable parameter space that can account for the CMS' $h \rightarrow \mu\tau$ signal interpretation. We also discuss how future searches for $\mu \rightarrow e\gamma$ and nuclear $\mu \rightarrow e$ conversion may offer further tests on the interactions of interest. Finally, we look at a few other processes that can be induced by the same operator. Especially, we find that the Z -boson decay $Z \rightarrow \mu\tau$ can have a branching ratio that is below its current empirical limit by merely less than an order of magnitude. We make our conclusions in section 4. An appendix contains some additional information and formulas.

2 Operators with minimal lepton-flavor violaton

In the SM plus three right-handed Majorana neutrinos, the renormalizable Lagrangian for lepton masses can be written as

$$\mathcal{L}_m = -(Y_\nu)_{kl} \bar{L}_{k,L} \nu_{l,R} \tilde{H} - (Y_e)_{kl} \bar{L}_{k,L} E_{l,R} H - \frac{1}{2} (M_\nu)_{kl} \overline{\nu_{k,R}^c} \nu_{l,R} + \text{H.c.}, \quad (2.1)$$

where $k, l = 1, 2, 3$ are implicitly summed over, $Y_{\nu,e}$ denote Yukawa coupling matrices, $L_{k,L}$ stands for left-handed lepton doublets, $\nu_{l,R}$ and $E_{l,R}$ represent right-handed neutrinos and charged leptons, respectively, $\tilde{H} = i\tau_2 H^*$ with τ_2 being the second Pauli matrix and H the Higgs doublet, M_ν is the Majorana mass matrix of $\nu_{l,R}$, and $\nu_{k,R}^c \equiv (\nu_{k,R})^c$, the

¹Various scenarios of leptonic MFV have been discussed in the literature [52–61].

superscript referring to charge conjugation. For the nonzero elements of M_ν taken to be much greater than those of $vY_\nu/\sqrt{2}$, the seesaw mechanism of type I is operational [62–70] and generates the light neutrinos’ mass matrix $m_\nu = -(v^2/2)Y_\nu M_\nu^{-1}Y_\nu^T = U_{\text{PMNS}}\hat{m}_\nu U_{\text{PMNS}}^T$, where $v \simeq 246$ GeV is the Higgs’s vacuum expectation value, U_{PMNS} denotes the Pontecorvo-Maki-Nakagawa-Sakata (PMNS [71, 72]) matrix, and $\hat{m}_\nu = \text{diag}(m_1, m_2, m_3)$ contains the light neutrinos’ eigenmasses. This suggests [73]

$$Y_\nu = \frac{i\sqrt{2}}{v}U_{\text{PMNS}}\hat{m}_\nu^{1/2}OM_\nu^{1/2}, \tag{2.2}$$

where O in general is a complex 3×3 matrix satisfying $OO^T = \mathbb{1}$, the right-hand side being a unit matrix, and can be parameterized as

$$O = e^{i\mathbf{R}}e^{\mathbf{R}'} , \quad \mathbf{R}^{(\prime)} = \begin{pmatrix} 0 & r_1^{(\prime)} & r_2^{(\prime)} \\ -r_1^{(\prime)} & 0 & r_3^{(\prime)} \\ -r_2^{(\prime)} & -r_3^{(\prime)} & 0 \end{pmatrix} \tag{2.3}$$

with $r_{1,2,3}$ and $r'_{1,2,3}$ being independent real constants. Hence nonvanishing $r'_{1,2,3}$ dictate how the Higgs couples to the right-handed neutrinos in a nontrivial way according to eq. (2.2). Hereafter, we concentrate on the possibility that the right-handed neutrinos are degenerate, so that $M_\nu = \mathcal{M}\mathbb{1}$. In this particular scenario, only the $e^{i\mathbf{R}}$ part of O matters physically [53].

The MFV hypothesis [50, 52] then implies that \mathcal{L}_m is formally invariant under the global flavor group $G_\ell = \text{SU}(3)_L \times \text{O}(3)_\nu \times \text{SU}(3)_E$. This entails that $L_{k,L}$, $\nu_{k,R}$, and $E_{k,R}$ belong to the fundamental representations of their respective flavor groups,

$$L_L \rightarrow V_L L_L, \quad \nu_R \rightarrow \mathcal{O}_\nu \nu_R, \quad E_R \rightarrow V_E E_R, \tag{2.4}$$

where $V_{L,E} \in \text{SU}(3)_{L,E}$ and $\mathcal{O}_\nu \in \text{O}(3)_\nu$ is an orthogonal real matrix [50, 52, 53]. Furthermore, under G_ℓ the Yukawa couplings transform in the spurion sense according to

$$Y_\nu \rightarrow V_L Y_\nu \mathcal{O}_\nu^T, \quad Y_e \rightarrow V_L Y_e V_E^\dagger. \tag{2.5}$$

Due to the symmetry under G_ℓ , we can work in the basis where $Y_e = \sqrt{2}\text{diag}(m_e, m_\mu, m_\tau)/v$ and the fields $\tilde{\nu}_{k,L}$, $\nu_{k,R}$, and E_k refer to the mass eigenstates. Explicitly, $(E_1, E_2, E_3) = (e, \mu, \tau)$. We can then express $L_{k,L}$ in relation to U_{PMNS} as

$$L_{k,L} = \begin{pmatrix} (U_{\text{PMNS}})_{kl}\tilde{\nu}_{l,L} \\ E_{k,L} \end{pmatrix}. \tag{2.6}$$

In the standard parametrization [51]

$$U_{\text{PMNS}} = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & e^{-i\delta}s_{13} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \text{diag}(e^{i\alpha_1/2}, e^{i\alpha_2/2}, 1), \tag{2.7}$$

where δ and $\alpha_{1,2}$ are the CP -violating Dirac and Majorana phases, respectively, $c_{kl} = \cos\theta_{kl}$, and $s_{kl} = \sin\theta_{kl}$.

To put together effective Lagrangians beyond the SM with MFV built-in, one inserts products of the Yukawa matrices among the pertinent fields to assemble G_ℓ -invariant operators that are singlet under the SM gauge group [50, 52]. Of interest here are the combinations

$$\mathbf{A} = Y_\nu Y_\nu^\dagger = \frac{2\mathcal{M}}{v^2} U_{\text{PMNS}} \hat{m}_\nu^{1/2} O O^\dagger \hat{m}_\nu^{1/2} U_{\text{PMNS}}^\dagger, \quad \mathbf{B} = Y_e Y_e^\dagger = \text{diag}(y_e^2, y_\mu^2, y_\tau^2), \quad (2.8)$$

where $y_f = \sqrt{2}m_f/v$. With these matrices, one can generally devise an object Δ as an infinite power series in them and their products, but it turns out to be resumable into only 17 terms [74, 75]. To maximize the new-physics effects, we assume that the right-handed neutrinos' mass \mathcal{M} is large enough to render the biggest eigenvalue of \mathbf{A} equal to unity, which conforms to the perturbativity requirement [59, 60, 74, 75]. Given that the eigenvalues of \mathbf{B} are at most $y_\tau^2 \sim 1 \times 10^{-4}$, we may consequently drop from Δ all the terms with \mathbf{B} , which would otherwise be needed in a study concerning CP violation [59–61]. Accordingly, the relevant building block is [61]

$$\Delta = \xi_1 \mathbb{1} + \xi_2 \mathbf{A} + \xi_4 \mathbf{A}^2, \quad (2.9)$$

where in our model-independent approach $\xi_{1,2,4}$ are free parameters expected to be at most of $\mathcal{O}(1)$, one or more of which could be suppressed or vanish, depending on the underlying theory. As $\text{Im}\xi_{1,2,4}$ are tiny [59, 60, 74, 75], we can further approximate $\Delta^\dagger = \Delta$.

One could then construct the desired G_ℓ -invariant effective Lagrangians that are SM gauge singlet. The one pertaining to $h \rightarrow \ell\ell'$ at tree level is given by [52]

$$\mathcal{L}_{\text{MFV}} = \frac{O_{RL}^{(e3)}}{\Lambda^2} + \text{H.c.}, \quad O_{RL}^{(e3)} = (\mathcal{D}^\rho H)^\dagger \bar{E}_R Y_e^\dagger \Delta \mathcal{D}_\rho L_L, \quad (2.10)$$

where the mass scale Λ characterizes the underlying heavy new-physics and the covariant derivatives $\mathcal{D}^\rho H = \partial^\rho H + i(g\tau_a W_a^\rho + g'B^\rho)H/2$ and $\mathcal{D}^\rho L = \partial^\rho L + i(g\tau_a W_a^\rho - g'B^\rho)L/2$ contain the usual $SU(2)_L \times U(1)_Y$ gauge fields W_a^ρ and B^ρ with coupling constants g and g' , respectively, and Pauli matrices τ_a , with summation over $a = 1, 2, 3$ being implicit. There are other dimension-six MFV operators involving H and leptons that have been written down [52],

$$\begin{aligned} i[H^\dagger \mathcal{D}_\rho H - (\mathcal{D}_\rho H)^\dagger H] \bar{L}_L \gamma^\rho \Delta_{LL} L_L, & \quad g' \bar{E}_R Y_e^\dagger \Delta_{RL} \sigma_{\rho\omega} H^\dagger L_L B^{\rho\omega}, \\ i[H^\dagger \tau_a \mathcal{D}_\rho H - (\mathcal{D}_\rho H)^\dagger \tau_a H] \bar{L}_L \gamma^\rho \Delta'_{LL} \tau_a L_L, & \quad g \bar{E}_R Y_e^\dagger \Delta'_{RL} \sigma_{\rho\omega} H^\dagger \tau_a L_L W_a^{\rho\omega}, \end{aligned} \quad (2.11)$$

with $\Delta_{LL,RL}^{(\prime)}$ being of the form of Δ in eq. (2.9) and having their own coefficients ξ_j , but these operators do not induce tree-level dilepton Higgs couplings. The same thing can be said of the comparatively more suppressed $i[H^\dagger \mathcal{D}_\rho H - (\mathcal{D}_\rho H)^\dagger H] \bar{E}_R \gamma^\rho Y_e^\dagger \Delta_{RR} Y_e E_R$. In the literature the operator $H^\dagger H \bar{E}_R Y_e^\dagger \Delta H^\dagger L_L$ is also often considered (e.g., [19]), but it can be shown to be related to $O_{RL}^{(e3)}$ and the other operators above. Explicitly, employing

the equations of motions for SM fields [76], one can derive [61]

$$\begin{aligned}
O_{RL}^{(e3)} + \text{H.c.} = & \frac{i}{8} [H^\dagger \mathcal{D}_\rho H - (\mathcal{D}_\rho H)^\dagger H] (\bar{L}_L \gamma^\rho \{\Delta, Y_e Y_e^\dagger\} L_L + 4 \bar{E}_R \gamma^\rho Y_e^\dagger \Delta Y_e E_R) \\
& + \frac{i}{8} [H^\dagger \tau_a \mathcal{D}_\rho H - (\mathcal{D}_\rho H)^\dagger \tau_a H] \bar{L}_L \gamma^\rho \{\Delta, Y_e Y_e^\dagger\} \tau_a L_L \\
& + \frac{i}{8} [H^\dagger \mathcal{D}_\rho H + (\mathcal{D}_\rho H)^\dagger H] \bar{L}_L \gamma^\rho [\Delta, Y_e Y_e^\dagger] L_L \\
& + \frac{i}{8} [H^\dagger \tau_a \mathcal{D}_\rho H + (\mathcal{D}_\rho H)^\dagger \tau_a H] \bar{L}_L \gamma^\rho [\Delta, Y_e Y_e^\dagger] \tau_a L_L \\
& + \frac{1}{8} [(4H^\dagger H/v^2 - 2)m_h^2 \bar{E}_R Y_e^\dagger \Delta H^\dagger L_L + 4 \bar{L}_L Y_e E_R \bar{E}_R Y_e^\dagger \Delta L_L \\
& \quad + \bar{E}_R Y_e^\dagger \Delta \sigma_{\rho\omega} H^\dagger (g' B^{\rho\omega} + g \tau_a W_a^{\rho\omega}) L_L + \text{H.c.}] \tag{2.12}
\end{aligned}$$

plus terms involving quark fields and total derivatives.² The third and fourth lines of this equation, which have $[\Delta, Y_e Y_e^\dagger]$, also supply contributions to $h \rightarrow \ell \ell'$, but they correspond to small, $\mathcal{O}(m_{\ell, \ell'}^2/m_h^2)$, effects that will be ignored later in eq. (3.3).

3 Decay amplitudes and numerical analysis

One can express the effective Lagrangian describing the Higgs decays $h \rightarrow \ell^- \ell'^+, \ell'^- \ell^+$ for $\ell \neq \ell'$ as

$$\mathcal{L}_{h\ell\ell'} = -\mathcal{Y}_{\ell\ell'} \bar{\ell} P_R \ell' - \mathcal{Y}_{\ell'\ell} \bar{\ell}' P_R \ell + \text{H.c.}, \tag{3.1}$$

where $\mathcal{Y}_{\ell\ell', \ell'\ell}$ denote the Yukawa couplings, which are in general complex. Hence the combined rate of $h \rightarrow \ell^- \ell'^+, \ell'^- \ell^+$ is

$$\Gamma_{h \rightarrow \ell\ell'} = \Gamma_{h \rightarrow \ell\bar{\ell}'} + \Gamma_{h \rightarrow \bar{\ell}\ell'} = \frac{m_h}{8\pi} (|\mathcal{Y}_{\ell\ell'}|^2 + |\mathcal{Y}_{\ell'\ell}|^2), \tag{3.2}$$

where the lepton masses have been neglected compared to m_h . The flavor-conserving decay $h \rightarrow \ell^- \ell^+$ has a rate of $\Gamma_{h \rightarrow \ell\bar{\ell}} = m_h |\mathcal{Y}_{\ell\ell}|^2 / (8\pi)$.

The MFV Lagrangian in eq. (2.10) contributes to both flavor-conserving and -violating Higgs decays. Including the SM part, we can write for $h \rightarrow E_k^- E_l^+$

$$\mathcal{Y}_{E_k E_l} = \delta_{kl} \mathcal{Y}_{E_k E_k}^{\text{SM}} - \frac{m_{E_l} m_h^2}{2\Lambda^2 v} \Delta_{kl}, \tag{3.3}$$

where $\mathcal{Y}_{E_k E_k}^{\text{SM}} = m_{E_k}/v$ at tree level. It follows that $|\mathcal{Y}_{\ell\ell'}| \ll |\mathcal{Y}_{\ell'\ell}|$ for $\ell\ell' = e\mu, e\tau, \mu\tau$ and $\mathcal{Y}_{\ell\ell}$ are real in our MFV scenario.

These couplings enter the amplitudes for a variety of lepton-flavor-violating processes, such as $\mu \rightarrow e\gamma$, via one- and two-loop diagrams. Therefore, they are subject to the pertinent empirical constraints [20–22], the most stringent of which we list here, assuming that the impact of these loop contributions is not much reduced by other new-physics

²The formula for $O_{RL}^{(e3)} + \text{H.c.}$ in the footnote 1 of ref. [61] has several terms missing and the wrong sign in the dipole ($\sigma_{\rho\omega}$) part. These errors have been corrected here in eq. (2.12).

effects. As we sketch in appendix A, the current bound $\mathcal{B}(\mu \rightarrow e\gamma)_{\text{exp}} < 5.7 \times 10^{-13}$ [51] translates into

$$\sqrt{|(\mathcal{Y}_{\mu\mu} + r_\mu)\mathcal{Y}_{\mu e} + 9.19\mathcal{Y}_{\mu\tau}\mathcal{Y}_{\tau e}|^2 + |(\mathcal{Y}_{\mu\mu} + r_\mu)\mathcal{Y}_{e\mu} + 9.19\mathcal{Y}_{e\tau}\mathcal{Y}_{\tau\mu}|^2} < 5.1 \times 10^{-7}, \quad (3.4)$$

where $r_\mu = 0.29$. From $\mathcal{B}(\tau \rightarrow e\gamma)_{\text{exp}} < 3.3 \times 10^{-8}$ [51], one extracts [19, 22]

$$|\mathcal{Y}_{\tau\tau} + r_\tau| \sqrt{|\mathcal{Y}_{\tau e}|^2 + |\mathcal{Y}_{e\tau}|^2} < 5.2 \times 10^{-4}, \quad (3.5)$$

where $r_\tau = 0.03$. In these inequalities, we have put more than two different couplings together, as they are generally affected by \mathcal{L}_{MFV} at the same time, and dropped smaller terms. The aforementioned CMS $h \rightarrow \mu\tau$ result under the no-signal assumption implies [8]

$$\sqrt{|\mathcal{Y}_{\tau\mu}|^2 + |\mathcal{Y}_{\mu\tau}|^2} < 3.6 \times 10^{-3}, \quad (3.6)$$

which is ~ 4 times stronger than the restraint [22] inferred from $\mathcal{B}(\tau \rightarrow \mu\gamma)_{\text{exp}} < 4.4 \times 10^{-8}$ [51] and encompasses the range

$$2.0 \times 10^{-3} < \sqrt{|\mathcal{Y}_{\tau\mu}|^2 + |\mathcal{Y}_{\mu\tau}|^2} < 3.3 \times 10^{-3} \quad (3.7)$$

implied by $\mathcal{B}(h \rightarrow \mu\tau) = (0.84_{-0.37}^{+0.39})\%$ in the CMS signal hypothesis [8].

The information on $h \rightarrow \mu^+\mu^-, \tau^+\tau^-$ recently acquired by ATLAS [4, 6] and CMS [5, 7] is also useful for restricting new physics in $\mathcal{Y}_{\mu\mu, \tau\tau}$. From the data described in section 1, we may require

$$|\mathcal{Y}_{\mu\mu}/\mathcal{Y}_{\mu\mu}^{\text{SM}}|^2 < 6.5, \quad 0.7 < |\mathcal{Y}_{\tau\tau}/\mathcal{Y}_{\tau\tau}^{\text{SM}}|^2 < 1.8, \quad (3.8)$$

where $\mathcal{Y}_{\mu\mu}^{\text{SM}} = 4.24 \times 10^{-4}$ and $\mathcal{Y}_{\tau\tau}^{\text{SM}} = 7.19 \times 10^{-3}$ in the SM from the rates $\Gamma_{h \rightarrow \mu\bar{\mu}}^{\text{SM}} = 894 \text{ eV}$ and $\Gamma_{h \rightarrow \tau\bar{\tau}}^{\text{SM}} = 257 \text{ keV}$ [77, 78] for $m_h = 125.1 \text{ GeV}$. These numbers allow one to see from eqs. (3.4) and (3.5), where r_μ and r_τ represent the 2-loop effects [19, 22], that the 2-loop contribution to $\mu \rightarrow e\gamma$ is dominant in constraining $\mathcal{Y}_{e\mu, \mu e}$, whereas the 1- and 2-loop effects on $\tau \rightarrow e\gamma$ are roughly comparable.

We now attempt to attain $|\mathcal{Y}_{\mu\tau}| \sim 0.003$ corresponding to the CMS hint of $h \rightarrow \mu\tau$ by scanning the coefficients $\xi_{1,2,4}$ in $\Delta = \xi_1 \mathbb{1} + \xi_2 \mathbf{A} + \xi_4 \mathbf{A}^2$ which enter the Yukawa couplings according to eq. (3.3) and consequently are subject to the restrictions in eqs. (3.4)–(3.6) and (3.8). Given that in our MFV scenario $\mathcal{Y}_{\ell\ell'} \propto m_{\ell'}$ if $\ell \neq \ell'$, from this point on we neglect $\mathcal{Y}_{\mu e, \tau e, \tau\mu}$ in comparison to $\mathcal{Y}_{e\mu, e\tau, \mu\tau}$, respectively.

Since A in eq. (2.8) can be realized in many different ways, we consider first the possibility that the orthogonal O matrix is real, in which case

$$\mathbf{A} = Y_\nu Y_\nu^\dagger = \frac{2\mathcal{M}}{v^2} U_{\text{PMNS}} \hat{m}_\nu U_{\text{PMNS}}^\dagger \quad (3.9)$$

and the right-handed neutrinos' Yukawa coupling matrix in eq. (2.2) simplifies to $Y_\nu \propto U_{\text{PMNS}} \hat{m}_\nu^{1/2}$, somewhat similar to its Dirac-neutrino counterpart [61]. Although U_{PMNS} has dependence on the Majorana phases $\alpha_{1,2}$, as in eq. (2.7), they drop out of eq. (3.9).

	$\frac{\alpha_1}{\pi}$	$\frac{\alpha_2}{\pi}$	r_1	r_2	r_3	$10^5 \xi_1/\Lambda^2$ (GeV ⁻²)	$10^5 \xi_2/\Lambda^2$ (GeV ⁻²)	$10^5 \xi_4/\Lambda^2$ (GeV ⁻²)	$\frac{\mathcal{Y}_{ee}}{\mathcal{Y}_{ee}^{\text{SM}}}$	$\frac{\mathcal{Y}_{\mu\mu}}{\mathcal{Y}_{\mu\mu}^{\text{SM}}}$	$\frac{\mathcal{Y}_{\tau\tau}}{\mathcal{Y}_{\tau\tau}^{\text{SM}}}$	$\frac{ \mathcal{Y}_{e\mu} }{10^{-6}}$	$\frac{ \mathcal{Y}_{e\tau} }{10^{-4}}$	$\frac{ \mathcal{Y}_{\mu\tau} }{10^{-3}}$
NH	0	0	0.81	-1.7	-0.89	-6.3	6.2	5.4	1.5	1.2	0.89	1.7	0.3	3.1
	0	0	-0.86	1.8	-0.92	-7.1	8.7	4.5	1.6	1.2	0.87	2.0	0.4	3.5
	0	0.23	0.74	-0.80	-0.20	4.9	-6.7	-5.9	0.63	0.93	1.3	1.7	2.2	3.2
IH	0	0	0.04	0.63	-0.93	-7.9	8.8	2.6	1.5	1.2	1.1	2.1	2.8	3.2
	0	0	0.02	-0.75	1.1	-5.7	3.8	8.1	1.4	1.1	0.90	2.4	1.3	3.3
	0.79	1.3	-0.61	-0.79	1.4	-5.3	5.0	7.6	1.4	1.0	0.84	1.2	0.4	3.5

Table 1. Higgs-lepton Yukawa couplings corresponding to sample values of the Majorana phases $\alpha_{1,2}$, the parameters $r_{1,2,3}$ of the complex O matrix, and the coefficients $\xi_{1,2,4}$ in the MFV building block Δ which can yield $|\mathcal{Y}_{\mu\tau}| \gtrsim 3 \times 10^{-3}$. The calculation of the NH (IH) results also relies on the measured neutrino mixing parameters in the case of normal (inverted) hierarchy of neutrino masses.

To proceed numerically, we employ the central values of neutrino mixing parameters from a recent fit to global neutrino data [79]. Most of the numbers depend on whether light neutrino masses have a normal hierarchy (NH), $m_1 < m_2 < m_3$, or an inverted one (IH), $m_3 < m_1 < m_2$. Since experimental information on the absolute scale of $m_{1,2,3}$ is still far from precise [51], for definiteness we select $m_1 = 0$ ($m_3 = 0$) in the NH (IH) case.

With the preceding choices, after exploring the $\xi_{1,2,4}$ parameter space, we find that $|\mathcal{Y}_{\mu\tau}|$ can only reach somewhere in the range of $(1-2) \times 10^{-4}$. This is caused by the constraint in eq. (3.4), without which the upper bound $|\mathcal{Y}_{\mu\tau}| < 0.0036$ could be easily saturated. Thus, to reproduce the signal range in eq. (3.7), the form of A in eq. (3.9) is not sufficient, and we instead need one with a less simple structure, to which we pay our attention next.³

A more promising possibility is that the O matrix in eq. (2.8) is complex, which leads to

$$A = Y_\nu Y_\nu^\dagger = \frac{2}{v^2} \mathcal{M} U_{\text{PMNS}} \hat{m}_\nu^{1/2} O O^\dagger \hat{m}_\nu^{1/2} U_{\text{PMNS}}^\dagger. \quad (3.10)$$

As mentioned in the previous section, one can express $O = e^{iR} e^{R'}$ with real antisymmetric matrices R and R' defined in eq. (2.3). Accordingly, we have

$$O O^\dagger = e^{2iR} = \mathbb{1} + iR \frac{\sinh(2\tilde{r})}{\tilde{r}} - 2R^2 \frac{\sinh^2 \tilde{r}}{\tilde{r}^2}, \quad \tilde{r} = \sqrt{r_1^2 + r_2^2 + r_3^2}, \quad (3.11)$$

and so nonzero $r_{1,2,3}$ can serve as extra free parameters that may allow us to achieve the desired size of $|\mathcal{Y}_{\mu\tau}|$. This can indeed be realized, as illustrated by the examples collected in table 1. The flavor-violating Yukawa couplings quoted in the last three columns have followed from their dependence on the elements of Δ determined using the listed sets of $\alpha_{1,2}$, $r_{1,2,3}$, and $\xi_{1,2,4}/\Lambda^2$ numbers, along with the central values of neutrino mixing parameters from ref. [79], again with $m_1 = 0$ ($m_3 = 0$) if the light neutrino masses have a normal (inverted) hierarchy. The table includes a couple of instances with nonvanishing Majorana phases $\alpha_{1,2}$, which are not yet measured and affect A , as $O O^\dagger$ in eq. (3.10) is not diagonal.

In the table, we also collect the corresponding flavor-conserving Yukawa couplings divided by their SM predictions, including \mathcal{Y}_{ee} for completeness, with $\mathcal{Y}_{ee}^{\text{SM}} = m_e/v =$

³A similar conclusion was drawn in ref. [44] from a semi-quantitative investigation focusing on an MFV contribution that corresponds to the ξ_2 term in our study.

2.08×10^{-6} . It is obvious that $\mathcal{Y}_{\ell\ell}$ can be altered sizeably with respect to their SM values. Therefore, measurements of $h \rightarrow \mu^+\mu^-, \tau^+\tau^-$ with improved precision in the future can offer complementary tests on the new contributions.

Based on our numerical exploration, there are a few more remarks we would like to make. First, we have noticed that the viable parameter ranges in the NH case are broader than their IH counterparts. Second, in many trials we observe that $|\mathcal{Y}_{e\tau}| \lesssim 0.1 |\mathcal{Y}_{\mu\tau}|$ for the hypothetical signal regions, as table 1 also shows. This pattern has implications that may be checked empirically in the future. Third, in the absence of either ξ_2 or ξ_4 the maximal $|\mathcal{Y}_{\mu\tau}|$ is somewhat lower than that when $\xi_{1,2,4}$ are all contributing, but at least some or all of the signal values in eq. (3.7) can be accommodated. However, if only ξ_2, ξ_4 , or $\xi_{2,4}$ are nonzero, $|\mathcal{Y}_{\mu\tau}|$ cannot exceed ~ 0.0018 .

Now, the six sample sets of parameter values in table 1 produce branching fractions of $\mu \rightarrow e\gamma$ and $\tau \rightarrow \mu\gamma$ in the ranges of $(1.4-5.4) \times 10^{-13}$ and $(1.6-2.0) \times 10^{-9}$, respectively, if other new-physics effects are negligible. The former numbers are within only a few times below the present bound $\mathcal{B}(\mu \rightarrow e\gamma)_{\text{exp}}$, whereas the latter are at least a factor of 20 less than $\mathcal{B}(\tau \rightarrow \mu\gamma)_{\text{exp}}$. They can be regarded as predictions testable by ongoing or future experiments looking for these decays if the CMS' indication of $h \rightarrow \mu\tau$ is substantiated by upcoming Higgs measurements and the signal range in eq. (3.7), or part of it, persists with increased data. Especially, the planned MEG II experiment on $\mu \rightarrow e\gamma$, with sensitivity expected to reach a few times 10^{-14} after 3 years of data taking [80, 81], will probe the above predictions for it.

As it turns out, if the forthcoming search for $\mu \rightarrow e\gamma$ still comes up empty, there could yet remain viable, but narrower, signal parameter regions. We illustrate this in table 2, assuming a possible future limit of $\mathcal{B}(\mu \rightarrow e\gamma) < 5 \times 10^{-14}$ [80], which amounts to replacing the right-hand side of eq. (3.4) with 1.5×10^{-7} , and also imposing the ratios $0.5 < \Gamma_{h \rightarrow \mu\bar{\mu}} / \Gamma_{h \rightarrow \mu\bar{\mu}}^{\text{SM}} < 1.5$ and $0.8 < \Gamma_{h \rightarrow \tau\bar{\tau}} / \Gamma_{h \rightarrow \tau\bar{\tau}}^{\text{SM}} < 1.2$ based on LHC Run-2 projections [82, 83]. Since the examples in table 2 yield $\mathcal{B}(\mu \rightarrow e\gamma) = (1.2-4.4) \times 10^{-14}$, they may be out of reach of MEG II, and so to probe them one will likely need to rely on experiments looking for nuclear $\mu \rightarrow e$ conversion, which promise a greater degree of sensitivity in the long run [81]. As discussed in appendix A, the existing data on $\mu \rightarrow e$ conversion in nuclei are not yet competitive to the current measured bound on $\mu \rightarrow e\gamma$ in constraining the Yukawa couplings. However, we also point out in the appendix that planned searches for $\mu \rightarrow e$ conversion, such as Mu2E and COMET [81], can be expected to test very well the parameter space represented by the examples in tables 1 and 2.

Finally, we discuss the contributions of \mathcal{L}_{MFV} in eq. (2.10) to some other processes. Expanding the operator, we have

$$\begin{aligned}
 O_{RL}^{(e3)} = & \frac{\Delta_{kl} m_{E_k}}{v} \bar{E}_k P_L \left(\partial_\eta E_l - ie A_\eta E_l + ig_L Z_\eta E_l + \frac{ig}{\sqrt{2}} W_\eta^- \nu_l \right) \partial^\eta h \\
 & + \frac{\Delta_{kl} g m_{E_k}}{v} \bar{E}_k P_L \left[\frac{i Z^\eta \partial_\eta E_l}{2c_w} - \frac{i W_\eta^- \partial^\eta \nu_l}{\sqrt{2}} + \left(\frac{eA \cdot Z}{2c_w} - \frac{g_L Z^2}{2c_w} + \frac{g}{2} W^+ \cdot W^- \right) E_l \right] (h+v),
 \end{aligned}
 \tag{3.12}$$

	$\frac{\alpha_1}{\pi}$	$\frac{\alpha_2}{\pi}$	r_1	r_2	r_3	$10^5 \xi_1 / \Lambda^2$ (GeV ⁻²)	$10^5 \xi_2 / \Lambda^2$ (GeV ⁻²)	$10^5 \xi_4 / \Lambda^2$ (GeV ⁻²)	$\frac{\mathcal{Y}_{ee}}{\mathcal{Y}_{ee}^{\text{SM}}}$	$\frac{\mathcal{Y}_{\mu\mu}}{\mathcal{Y}_{\mu\mu}^{\text{SM}}}$	$\frac{\mathcal{Y}_{\tau\tau}}{\mathcal{Y}_{\tau\tau}^{\text{SM}}}$	$\frac{ \mathcal{Y}_{e\mu} }{10^{-6}}$	$\frac{ \mathcal{Y}_{e\tau} }{10^{-4}}$	$\frac{ \mathcal{Y}_{\mu\tau} }{10^{-3}}$
NH	0	0	-0.53	0.73	-0.40	6.0	-0.7	-9.5	0.53	0.79	1.1	0.6	0.2	2.7
	0	0.4	0.68	-0.80	-0.15	-5.4	-2.3	12	1.4	1.2	0.93	0.3	0.5	2.6
IH	0	0	0.0	-0.73	1.1	-4.7	-1.9	11	1.4	1.1	0.96	0.5	0.1	2.5
	0.8	1.3	-0.60	-0.81	1.4	-6.5	9.4	1.1	1.5	1.2	1.0	0.1	0.5	2.9

Table 2. The same as table 1, except the $\mu \rightarrow e\gamma$ and $h \rightarrow \mu\bar{\mu}, \tau\bar{\tau}$ constraints are replaced with their projected future experimental limits, as described in the text.

where $g_L = g(s_w^2 - 1/2)/c_w$ and $c_w = \sqrt{1 - s_w^2} = gv/(2m_Z) = m_W/m_Z$. Evidently, \mathcal{L}_{MFV} not only induces the already addressed $h \rightarrow \ell\bar{\ell}'$ couplings, but also contributes to the two-body decays of the weak bosons, $Z \rightarrow \ell\bar{\ell}'$ and $W \rightarrow \tau\nu_l$, as well as to three- and four-body modes, such as $h \rightarrow \ell\bar{\ell}'\gamma, \nu\ell W^+, \ell\bar{\ell}'\gamma Z$. Since the latter are more suppressed by phase space, we deal with only the two-body Z and W decays. The other operators in eq. (2.11) can also affect $Z \rightarrow \ell\bar{\ell}'$ and $W \rightarrow \tau\nu_l$, but here we entertain the possibility that their impact is comparatively unimportant. Accordingly, from eq. (3.12) we derive

$$\begin{aligned}
 \mathcal{M}_{Z \rightarrow E_k \bar{E}_l} &= \bar{u}_{E_k} \left[\delta_{kl} \not{\epsilon}_Z (g_L P_L + g_R P_R) + \frac{\Delta_{kl} m_Z}{\Lambda^2 v} \left(m_{E_k} P_L \epsilon_Z \cdot p_{E_l} - m_{E_l} P_R \epsilon_Z \cdot p_{E_k} \right) \right] v_{E_l}, \\
 \mathcal{M}_{W \rightarrow \tau \nu_l} &= \bar{u}_\tau \left(\frac{\delta_{3l} g}{\sqrt{2}} \not{\epsilon}_W + \frac{\sqrt{2} \Delta_{3l} m_\tau m_W}{\Lambda^2 v} \epsilon_W \cdot p_\tau \right) P_L v_{\nu_l}.
 \end{aligned} \tag{3.13}$$

where $g_R = gs_w^2/c_w$ and we have included the SM terms in these amplitudes. Hence, neglecting lepton masses compared to m_Z , we arrive at

$$\Gamma_{Z \rightarrow \mu \bar{e}} = \Gamma_{Z \rightarrow \mu \bar{e}} \simeq \frac{|\Delta_{12} m_\mu|^2 m_Z^5}{192 \Lambda^4 \pi v^2} = \frac{|\mathcal{Y}_{e\mu}|^2 m_Z^5}{48 \pi m_h^4} \tag{3.14}$$

and similarly for $Z \rightarrow e\tau, \mu\tau$. Thus, for, say, $|\mathcal{Y}_{e\mu}| = 2.1 \times 10^{-6}$, $|\mathcal{Y}_{e\tau}| = 2.8 \times 10^{-4}$, and $|\mathcal{Y}_{\mu\tau}| = 0.0032$ from table 1, we get

$$\mathcal{B}(Z \rightarrow e^\pm \mu^\mp) = 6.0 \times 10^{-13}, \quad \mathcal{B}(Z \rightarrow e^\pm \tau^\mp) = 1.1 \times 10^{-8}, \quad \mathcal{B}(Z \rightarrow \mu^\pm \tau^\mp) = 1.4 \times 10^{-6}. \tag{3.15}$$

For comparison, the experimental limits are [51]

$$\begin{aligned}
 \mathcal{B}(Z \rightarrow e^\pm \mu^\mp)_{\text{exp}} &< 1.7 \times 10^{-6}, & \mathcal{B}(Z \rightarrow e^\pm \tau^\mp)_{\text{exp}} &< 9.8 \times 10^{-6}, \\
 \mathcal{B}(Z \rightarrow \mu^\pm \tau^\mp)_{\text{exp}} &< 1.2 \times 10^{-5}
 \end{aligned} \tag{3.16}$$

at 95% CL. We see that the predicted $\mathcal{B}(Z \rightarrow \mu\tau)$ is below its experimental bound by only less than a factor of 10. Therefore, $Z \rightarrow \mu\tau$ is potentially more testable than $Z \rightarrow e\mu, e\tau$, and the quest for it can provide a complementary check on \mathcal{L}_{MFV} .

Neglecting lepton masses compared to $m_{W,Z}$, we also obtain from eq. (3.13)

$$\begin{aligned}\Gamma_{Z \rightarrow E_k \bar{E}_k} &= \frac{m_Z}{24\pi} \left(g_L^2 + g_R^2 + \frac{\Delta_{kk}^2 m_{E_k}^2 m_Z^4}{4\Lambda^4 v^2} \right), \\ \Gamma_{W \rightarrow \tau \nu} &= \frac{m_W}{48\pi} \left(g^2 + \frac{\Delta_{33}^2 m_\tau^2 m_W^4}{2\Lambda^4 v^2} \right) + \frac{(|\Delta_{31}|^2 + |\Delta_{32}|^2) m_\tau^2 m_W^4}{96\Lambda^4 \pi v^2},\end{aligned}\quad (3.17)$$

where in the $W \rightarrow \tau \nu$ formula we have summed over the 3 neutrino flavors. For the parameter values in table 1, the nonstandard terms in $\Gamma_{Z \rightarrow E_k \bar{E}_k}$ and $\Gamma_{W \rightarrow \tau \nu}$ are tiny, being smaller than the SM parts by more than 4 orders of magnitude.

Before ending this section, we would like to note that all the preceding analysis can be repeated within the context of the type-III seesaw model [84] with MFV, which is very similar to the type-I case addressed in this study if the triplet leptons in the former are as heavy as the right-handed neutrinos in the latter [61]. However, in the type-II seesaw model [85–88] with MFV, the Yukawa coupling matrix of the triplet scalars does not possess the special feature that Y_ν has with regard to the O matrix [61] that allows $\mathcal{Y}_{\mu\tau}$ to become large enough to explain the CMS $h \rightarrow \mu\tau$ signal hypothesis.

4 Conclusions

We have explored the possibility that the slight excess of $h \rightarrow \mu\tau$ events recently detected in the CMS experiment has a new-physics origin. Adopting in particular the effective theory framework of MFV, we consider the SM extended with the type-I seesaw mechanism and an effective dimension-six operator responsible for the flavor-violating dilepton Higgs decay. We demonstrate that to account for the tentative $h \rightarrow \mu\tau$ signal, with a branching fraction of order 1%, the Yukawa coupling matrix of the right-handed neutrinos needs to have a nontrivial structure because of the stringent empirical constraints. To illustrate this, we present several benchmark points that have survived the restrictions from the existing $\mu \rightarrow e\gamma$, $\tau \rightarrow e\gamma$, and $h \rightarrow \mu\bar{\mu}, \tau\bar{\tau}$ data. The viable parameter space can be probed further by upcoming LHC measurements and future quests for charged-lepton-flavor violation. Lastly, we examine a few other transitions that arise from the same dimension-six operator, among which $Z \rightarrow \mu\tau$ can have a predicted branching ratio merely less than 10 times below its current empirical limit and hence potentially also testable in near-future searches.

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A Constraints from $\mu \rightarrow e\gamma$ decay and $\mu \rightarrow e$ conversion

The effective Lagrangian for $\mu \rightarrow e\gamma$ can be expressed as

$$\mathcal{L}_{\mu \rightarrow e\gamma} = \frac{\sqrt{\alpha\pi} m_\mu}{4\pi^2} \bar{e} \sigma^{\rho\omega} (\mathcal{C}_L P_L + \mathcal{C}_R P_R) \mu F_{\rho\omega}, \quad (A.1)$$

where $\alpha = 1/137$ is the fine structure constant, $P_{L,R} = (1 \mp \gamma_5)/2$, and $F_{\rho\omega}$ is the electromagnetic field strength tensor. This leads to the decay rate

$$\Gamma_{\mu \rightarrow e\gamma} = \frac{\alpha m_\mu^5}{64\pi^4} (|\mathcal{C}_L|^2 + |\mathcal{C}_R|^2), \quad (\text{A.2})$$

The Wilson coefficients $\mathcal{C}_{L,R}$ receive contributions from Higgs-mediated one-loop and two-loop [89] diagrams, $\mathcal{C}_{L,R} = \mathcal{C}_{L,R}^{\text{1loop}} + \mathcal{C}_{L,R}^{\text{2loop}}$. Given that $\mathcal{Y}_{\ell\ell}$ is real and $|\mathcal{Y}_{ee}| \ll |\mathcal{Y}_{\mu\mu}|$, one finds [22]

$$\begin{aligned} \mathcal{C}_R^{\text{1loop}} &\simeq \frac{\mathcal{Y}_{\mu\mu}\mathcal{Y}_{e\mu}}{2m_h^2} \left(\log \frac{m_h}{m_\mu} - \frac{2}{3} \right) + \frac{m_\tau \mathcal{Y}_{e\tau} \mathcal{Y}_{\tau\mu}}{2m_\mu m_h^2} \left(\log \frac{m_h}{m_\tau} - \frac{3}{4} \right), \\ \mathcal{C}_R^{\text{2loop}} &\simeq \frac{0.055 m_\tau \mathcal{Y}_{e\mu}}{m_\mu m_h^2} \end{aligned} \quad (\text{A.3})$$

and $\mathcal{C}_L^{\text{1loop,2loop}}$ obtainable from $\mathcal{C}_R^{\text{1loop,2loop}}$ with the replacements $\mathcal{Y}_{\ell\ell'} \rightarrow \mathcal{Y}_{\ell'\ell}^*$. Here we suppose that there are no other new-physics contributions that can bring about destructive interference with these coefficients. Thus, putting together these formulas with the latest experimental bound [51] $\mathcal{B}(\mu \rightarrow e\gamma)_{\text{exp}} < 5.7 \times 10^{-13}$, we arrive at eq. (3.4) for $m_h = 125.1$ GeV, which is consistent with the most recent measurement [90].

The effective Lagrangian for $\mu \rightarrow e$ conversion in nuclei is [91]

$$\mathcal{L}_{\mu \rightarrow e} = \frac{\sqrt{\alpha\pi} m_\mu}{4\pi^2} \bar{e} \sigma^{\rho\omega} (\mathcal{C}_L P_L + \mathcal{C}_R P_R) \mu F_{\rho\omega} - \frac{1}{2} \sum_q \bar{e} (g_{LS}^q P_R + g_{RS}^q P_L) \mu \bar{q} q, \quad (\text{A.4})$$

where q runs over all quark flavors, we have displayed only the most important terms for our purposes, and, if $\mathcal{Y}_{\ell\ell'}$ are the only LFV sources, $\mathcal{C}_{L,R}$ are already written down in the preceding paragraph and [22]

$$g_{LS}^q = \frac{-2m_q \mathcal{Y}_{\mu e}^*}{m_h^2 v}, \quad g_{RS}^q = \frac{-2m_q \mathcal{Y}_{e\mu}}{m_h^2 v}. \quad (\text{A.5})$$

The $\mu \rightarrow e$ conversion rate in nucleus \mathcal{N} is then given by [91]

$$\mathcal{B}(\mu\mathcal{N} \rightarrow e\mathcal{N}) = \frac{m_\mu^5}{\omega_{\text{capt}}^{\mathcal{N}}} \left| \frac{\sqrt{\alpha\pi} \mathcal{C}_L D_{\mathcal{N}}}{8\pi^2} - \tilde{g}_{LS}^{(p)} S_{\mathcal{N}}^{(p)} - \tilde{g}_{LS}^{(n)} S_{\mathcal{N}}^{(n)} \right|^2 + (L \rightarrow R), \quad (\text{A.6})$$

$$\tilde{g}_{LS}^{(N)} = \sum_q \frac{g_{LS}^q}{m_q} f_q^{(N)} m_N = \frac{-2m_N \mathcal{Y}_{\mu e}^*}{m_h^2 v} \sum_q f_q^{(N)}, \quad f_q^{(N)} = \frac{\langle N | m_q \bar{q} q | N \rangle}{m_N}, \quad N = p, n, \quad (\text{A.7})$$

where $D_{\mathcal{N}}$ and $S_{\mathcal{N}}^{(p,n)}$ are dimensionless integrals representing the overlap of electron and muon wave functions for \mathcal{N} and $\omega_{\text{capt}}^{\mathcal{N}}$ is the rate of muon capture in \mathcal{N} . Based on the current experimental limits on $\mu \rightarrow e$ transition in various nuclei [51, 92] and the corresponding overlap integral and $\omega_{\text{capt}}^{\mathcal{N}}$ values [91], one expects that the $\mathcal{N} = \text{Au}$ and Ti data may supply the most consequential restrictions. The evaluation of $\mathcal{B}(\mu\mathcal{N} \rightarrow e\mathcal{N})$ for these two nuclei, respectively, requires $D_{\text{Au}} = 0.189$, $D_{\text{Ti}} = 0.087$, $S_{\text{Au}}^{(p)} = 0.0614$, $S_{\text{Au}}^{(n)} = 0.0918$, $S_{\text{Ti}}^{(p)} = 0.0368$, $S_{\text{Ti}}^{(n)} = 0.0435$, $\omega_{\text{capt}}^{\text{Au}} = 13.07 \times 10^6/\text{s}$, and $\omega_{\text{capt}}^{\text{Ti}} = 2.59 \times 10^6/\text{s}$ [91], as

well as the latest determination of the sum of the nucleon matrix elements, $\Sigma_q f_q^{(p,n)} = 0.305 \pm 0.009$ [93],⁴ which lies around the lower end of the ranges from some of earlier estimates [95, 96].

If we impose the measured bound $\mathcal{B}(\mu\text{Au} \rightarrow e\text{Au})_{\text{exp}} < 7 \times 10^{-13}$ [51], instead of eq. (3.4), but still apply eqs. (3.5), (3.6), and (3.8), we end up with $|\mathcal{Y}_{e\mu}| < 1.6 \times 10^{-5}$, which is compatible with the finding of ref. [97]. If we use $\mathcal{N} = \text{Ti}$ with $\mathcal{B}(\mu\text{Ti} \rightarrow e\text{Ti})_{\text{exp}} < 6.1 \times 10^{-13}$ [92], instead of $\mathcal{N} = \text{Au}$, we get the somewhat stricter $|\mathcal{Y}_{e\mu}| < 1.3 \times 10^{-5}$. These limitations are roughly 5 to 13 times higher than the range of results $|\mathcal{Y}_{e\mu}| = (1.2\text{--}2.4) \times 10^{-6}$ quoted in table 1, demonstrating that the present data on nuclear $\mu \rightarrow e$ conversion are not yet competitive to $\mathcal{B}(\mu \rightarrow e\gamma)_{\text{exp}}$ in restricting especially $\mathcal{Y}_{e\mu}$, which is also known in the literature [20–22, 97]. Nevertheless, the leading planned searches for $\mu \rightarrow e$ conversion, Mu2E and COMET, which utilize aluminum as the target material [81], will likely be able to probe the parameter space represented by the examples in both tables 1 and 2. More precisely, from the sets of sample numbers in these tables, together with the aluminum parameters $D_{\text{Al}} = 0.0362$, $S_{\text{Al}}^{(p)} = 0.0155$, $S_{\text{Al}}^{(n)} = 0.0167$, and $\omega_{\text{capt}}^{\text{Al}} = 0.7054 \times 10^6/\text{s}$ [91], we obtain $\mathcal{B}(\mu\text{Al} \rightarrow e\text{Al}) = (0.1\text{--}9.0) \times 10^{-15}$, which are within reach of Mu2E and COMET, expected to have sensitivity levels under 10^{-16} or better after several years of running [81].

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⁴This sum partly depends on the pion-nucleon σ -term determined in ref. [93] which agrees with that previously calculated in ref. [94].

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