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Reconciling leptogenesis with observable $\mu \longrightarrow e\gamma$ rates

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Abstract: We perform a detailed analysis of thermal leptogenesis in the framework of seesaw models which approximately conserve lepton number. These models are known to allow for large Yukawa couplings and a low seesaw scale in agreement with neutrino mass constraints, and hence to lead to large lepton flavour violating rates that can be probed experimentally. Although large Yukawa couplings lead to (inverse) decay rates much larger than the Hubble expansion rate, we show that the leptogenesis washout induced is generically small if the mass splitting between the right-handed neutrinos is small enough. As a result, large lepton flavour violating rates are compatible with successful leptogenesis. We emphasize that this scenario does not require any particular flavour structure. A small splitting is natural and radiatively stable in this context because it is protected by the lepton number symmetry.

Keywords: Beyond Standard Model, Cosmology of Theories beyond the SM, Neutrino Physics

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Contents

1 Introduction

The seesaw models provide an attractive and straightforward explanation for both the recently observed tiny neutrino masses and the baryon asymmetry of the universe [1]. However, it is a very difficult task to probe these models experimentally, especially the most studied "type-I" seesaw model [2–6] based on the existence of gauge singlet right-handed (RH) neutrinos. In the generic situation, where *lepton number* violation and *lepton flavour* violation are generated at the same scale, neutrino mass constraints require either an experimentally unreachable heavy scale for the mass of the RH neutrinos, or tiny Yukawa couplings (leading to suppressed RH neutrino production cross sections). In both cases the rate of rare lepton processes, such as $\mu \to e\gamma$, is expected to be highly suppressed.

Without assuming any extra new physics beyond the right-handed neutrinos,[1] there exists however a theoretically motivated class of seesaw models where tiny neutrino masses do not imply that the lepton flavour violation is suppressed: the seesaw models which approximately conserve the total lepton number L [9]–[16]. These models are based on the fact that, even in the presence of both RH neutrino masses and neutrino Yukawa couplings, it is possible to conserve L, and thus to keep the left-handed neutrinos massless. Hence the lepton flavour violating scale can be disconnected from the lepton number violating one (see [16] for a detailed discussion). As a result, the Yukawa couplings are not constrained to be small even if the RH neutrino mass scale is low. Large lepton-flavour-violating (LFV) –but L-conserving– processes can be induced, with rates that could be

[1] In the following we assume a non-supersymmetric setup. In supersymmetric seesaw models it is known that large flavour violating rates can be induced in agreement with leptogenesis, see e.g. [7, 8].

observed if the RH neutrino mass scale is sufficiently low, a few tens of TeV or less. Neutrino masses can be explained subsequently by introducing small L-violating perturbations in the Yukawa coupling matrix and/or the RH neutrino mass matrix.

Successful baryogenesis via leptogenesis is known to be feasible at the TeV scale [17, 18] through the resonant leptogenesis mechanism [17, 19, 20], which, through the virtual N propagator of the one-loop self-energy diagram [21–23], requires a quasi-degeneracy of at least two of the RH neutrinos. A nice feature of the approximately L-conserving models is that they precisely yield a quasi-degenerate spectrum, with a mass splitting proportional to the small L-violating entries. Successful leptogenesis is achievable in this case, at least for Yukawa couplings suppressed enough [24].

In this letter we investigate the following question: is successful leptogenesis compatible with observable rare lepton processes? In ref. [25] it was shown to be feasible considering a particular flavour structure (suppressed τ Yukawa couplings) and with an extra SO(3) symmetry between the 3 RH neutrinos. In the following we show that it is generically possible even with two RH neutrinos and without the need to assume any particular flavour structure. This is due to the fact that the large Yukawa couplings which can lead to observable rare lepton processes are precisely the ones which do not break lepton number, therefore not necessarily causing any washout of the lepton asymmetry.

In section 2 we introduce the model with approximate L conservation which is going to be at the center of our study. In section 3 we show explicitly how the washout is suppressed due to subtle interference effects in the L-conserving limit. In section 4 we compute the *CP* asymmetry and, in section 5, study the consequences of our finding for the observability of rare lepton processes. In section 6 we conclude and emphasize the generality of our results.

2 The model

We consider for illustration the simple Type A model considered in [16] with only two RH neutrinos, $N_{1,2}$ (see also [24]). With the charge assignment $L=1$, $L=-1$ to N_1 and N_2 respectively, the Lagrangian can be split into a L-conserving part and a L-violating one:

$$\mathcal{L}_L = \mathrm{i}\overline{N}_i \not{\partial} N_i - \left(Y_\alpha \overline{N}_1 \tilde{\phi}^\dagger \ell_{L\alpha} + \frac{1}{2} M \overline{N}_1 N_2^c + h.c. \right) , \tag{2.1}$$

$$\mathcal{L}_{\not{L}} = - \left(Y'_\alpha \overline{N}_2 \tilde{\phi}^\dagger \ell_{L\alpha} + \frac{1}{2}\mu_1 \overline{N}_1 N_1^c + \frac{1}{2}\mu_2 \mathrm{e}^{\mathrm{i}\alpha} \overline{N}_2 N_2^c + h.c. \right) . \tag{2.2}$$

with $\phi \equiv (\phi^+ \phi^0)^T$ the ordinary Higgs doublet. Note that in this Lagrangian μ_1, μ_2 and M are real parameters. The phases of μ_1 and M have been rotated away. The Yukawa couplings Y and the RH neutrino mass M do not break lepton number, i.e. do not induce any L-violating dimension 5 operator for neutrino masses. But they do generate a dimension 6 L-conserving operator, and hence rare lepton processes, which can be large. For example a $\mu \to e\gamma$ rate of order the experimental upper bound $\mathrm{Br}(\mu \to e\gamma) \simeq 10^{-11}$ can be obtained for $M \simeq 1\,\mathrm{TeV}$ if the Y couplings are of order few 10^{-2}. On the other hand, the couplings Y', as well as the diagonal mass entries μ_1 and μ_2, do break lepton number and are therefore constrained to be small by the neutrino masses: $Y' \ll Y$ and μ_1, $\mu_2 \ll M$. This can be seen by diagonalizing the neutral state mass matrix which results from eqs. (2.1)–(2.2)

$$M_\nu = \begin{pmatrix} 0 & Y_\alpha^T\, v & Y_\alpha'^T\, v \\ Y_\alpha\, v & \mu_1 & M \\ Y'_\alpha\, v & M & \mu_2 e^{i\alpha} \end{pmatrix}, \tag{2.3}$$

and which, at first order in μ_1, μ_2 and Y', gives the following light neutrino mass matrix (with $v = 174\,\mathrm{GeV}$):

$$m_\nu = v^2 \left(Y'^T \frac{1}{M} Y + Y^T \frac{1}{M} Y' \right) - v^2 \left(Y^T \frac{1}{M} \mu_2 e^{i\alpha} \frac{1}{M} Y \right). \tag{2.4}$$

For $M \sim 1\,\mathrm{TeV}$ a neutrino mass of order $0.1\,\mathrm{eV}$ typically requires $Y'Y \sim 10^{-12}$ and/or $\mu_2 Y^2/M \sim 10^{-12}$. For $\mu_2 = 0$ this model has also the nice and rather unique property to allow for the full reconstruction of the flavour structure of the model from the values of the light neutrino mass matrix entries [16]. For $\mu_2 \neq 0$ it is not the case but still the full flavour structure of the L-conserving rare lepton processes can be reconstructed, leading to definite predictions for the various rare lepton processes, up to an overall normalization.

To consider leptogenesis in this framework, it is convenient to go to the basis where the RH neutrino mass matrix is diagonal with real and positive entries. The Lagrangian is then given by

$$\mathcal{L} = i\overline{N_i}\, \not\partial N_i - \left(h_{i\alpha}\, \overline{N_i} \tilde{\phi}^\dagger\, \ell_{L\alpha} + \frac{1}{2} M_i \overline{N_i} N_i^c + h.c. \right), \quad (i = 1,2;\ \alpha = e, \mu, \tau). \tag{2.5}$$

To first order in μ/M and Y', the mass eigenvalues are given by

$$M_{1,2} \simeq M \mp \frac{1}{2}\mu, \tag{2.6}$$

where we defined $\mu_1 + \mu_2 e^{i\alpha} \equiv \mu e^{i\phi}$, so that $\mu^2 = \mu_1^2 + \mu_2^2 + 2\mu_1\mu_2 \cos\alpha$. As for the Yukawa couplings, again to first order in μ/M and Y', we find

$$h_{1\alpha} \simeq \frac{i}{\sqrt{2}} e^{-i(\phi-\lambda)/2} \left[\left(1 + \frac{\mu_2^2 - \mu_1^2}{4M\mu} \right) e^{i\phi} Y_\alpha - Y'_\alpha \right], \tag{2.7}$$

$$h_{2\alpha} \simeq \frac{1}{\sqrt{2}} e^{-i(\phi+\lambda)/2} \left[\left(1 - \frac{\mu_2^2 - \mu_1^2}{4M\mu} \right) e^{i\phi} Y_\alpha + Y'_\alpha \right], \tag{2.8}$$

where

$$\lambda = \sin\alpha \frac{\mu_1 \mu_2}{\mu M}. \tag{2.9}$$

3 Washout from $\Delta L = 2$ scatterings

The RH neutrinos being unstable particles, one cannot in principle define asymptotic free states for them. However, as shown in [17, 26], their interaction properties (decay, inverse decay...) can be deduced from the $\Delta L = 2$ scatterings mediated by on-shell heavy neutrinos $\ell_\alpha \tilde{\phi}^\dagger \leftrightarrow \bar{\ell}_\beta \tilde{\phi}$. The corresponding $\gamma^{\rm tot}_{\Delta L=2,\alpha}$ interaction rate enters in the Boltzmann equation for the lepton flavour α as

$$s\, z\, H(z) \frac{dY_{\ell_\alpha}}{dz} = \sum_{i=1}^{2} \epsilon_{i\alpha} \left(\frac{Y_{N_i}(z)}{Y^{\rm eq}_{N_i}(z)} - 1 \right) \gamma^D_{N_i} - \sum_{i=1}^{2} \frac{Y_{\ell_\alpha}(z)}{Y^{\rm eq}_{\ell}(z)}\, 2\, \gamma^{\rm tot}_{\Delta L=2,i\alpha}\,, \tag{3.1}$$

where s is the entropy density, $z \equiv M_1/T$, $H(z)$ is the Hubble expansion rate, $Y_X \equiv n_X/s$ is the number density of particle X normalized to the entropy density and $\epsilon_{i\alpha}$ is the *CP* asymmetry from the RH neutrino N_i for the flavour α. Note that we consider a Boltzmann equation for each flavour, even though our results will not require to assume any particular flavour pattern.

One can separate the $\Delta L = 2$ scattering into an on-shell and off-shell contribution, $\gamma^{\rm tot}_{\Delta L=2,\alpha} = \gamma^{\rm on}_{\Delta L=2,\alpha} + \gamma^{\rm off}_{\Delta L=2,\alpha}$. As shown in the appendix of [27], in usual leptogenesis scenarios, the on-shell part of the $\Delta L = 2$ scattering, which is the dominant contribution, yields the inverse decays of leptons into heavy neutrinos,

$$\gamma^{\rm on}_{\Delta L=2,i\alpha} = \frac{\gamma^D_{N_i,\alpha}}{4} \equiv \frac{1}{4} n^{\rm eq}_{N_i} \frac{\mathcal{K}_1}{\mathcal{K}_2} \Gamma_{N_i,\alpha} = \frac{4}{z} \frac{M_i^4}{128\pi^3} |h_{i\alpha}|^2 \mathcal{K}_1(z\, M_i/M_1) \tag{3.2}$$

where $\mathcal{K}_{1,2}(z)$ are the modified Bessel functions. Note that $\sum_\alpha \gamma^D_{N_i,\alpha} = \gamma^D_{N_i}$. Inverse decays contribute to a depletion of the lepton asymmetry, parametrized by the washout parameter $K_{i\alpha} \equiv \Gamma_{N_i,\alpha}/H(T = M_{N_i})$, where the N_i decay width is $\Gamma_{N_i,\alpha} = |h_{i\alpha}|^2\, M_i/8\,\pi$. With quasi-degenerate RH neutrinos the relevant washout parameter is given by the sum $K_\alpha \equiv K_{1\alpha} + K_{2\alpha} = |Y_\alpha|^2\, M/8\,\pi/H(M) + \mathcal{O}(Y'^2, \mu^2/M^2)$.

There is an apparent contradiction here. On the one hand, it is clear that in the L-conserving limit $\Delta L = 2$ scatterings must vanish. On the other hand, the on-shell contribution quoted above clearly does not vanish, since the decay width is dominated by the L-conserving Y couplings, $|h|^2 \simeq |Y|^2/2$. For example, if the Y couplings are large, say 5×10^{-2}, and if the RH neutrino mass M is as low as 1 TeV, the washout parameter K_α is huge, about 10^{11}, which gives a baryon asymmetry much too small. However, this huge washout is independent of the L-violating parameters, i.e. it would be present even in the absence of L-violation, which is clearly erroneous.

In the following we show that this apparent contradiction disappears when one takes into account, on top of the inverse decay terms, the usually neglected interference terms involving different N_i's. To this end we start from the $\Delta L = 2$ scattering reaction density as given from the reduced cross section $\hat{\sigma}_{\Delta L=2,\alpha}$:

$$\gamma^{\rm tot}_{\Delta L=2,\alpha}(z) = \frac{1}{z} \frac{M_1^4}{64\,\pi^4} \int_{x_{\rm thr}}^{\infty} dx\, \sqrt{x}\, \hat{\sigma}_{\Delta L=2,\alpha}(x)\, \mathcal{K}_1(z\,\sqrt{x})\,, \tag{3.3}$$

where $x \equiv s/M_1^2$, and with the scattering threshold given by $x_{\rm thr} = (m_h + m_\ell)^2/M^2 \simeq 0$.

The s-channel reduced cross section is given by[2]

$$\hat{\sigma}_{\Delta L=2,\alpha} = |h_{1\alpha}|^2 A_{11} \sum_\beta |h_{1\beta}|^2 + |h_{2\alpha}|^2 A_{22} \sum_\beta |h_{2\beta}|^2 + 2\,\mathrm{Re}\left[(h_{1\alpha} h^\star_{2\alpha})\, A_{12} \sum_\beta \left(h_{1\beta} h^\star_{2\beta}\right)\right], \tag{3.4}$$

where the $h_{\alpha i}$'s are the neutrino Yukawa couplings of eqs. (2.7)–(2.8). In the above expression

$$A_{ij} = \frac{x\,\sqrt{a_i a_j}}{4\pi P_i P_j^\star}\,, \tag{3.5}$$

with $P_i^{-1} = \left(x - a_i + i\sqrt{a_i c_i}\right)^{-1}$ the dressed inverse propagators, whose poles are located at $\hat{a}_i = a_i - i\sqrt{a_i c_i}$, with $a_i = (M_i/M_1)^2$ and $c_i = (\Gamma_i/M_1)^2$. In the following, whenever the small L-violating $c_2 - c_1$ difference is irrelevant we will take $c_1 = c_2 = c$.

Using a narrow width approximation,[3] the $\Delta L = 2$ washout term is given by

$$\begin{aligned}\gamma^{\mathrm{on}}_{\Delta L=2,\alpha} = {} & \frac{1}{z}\frac{M_1^4}{64\,\pi^4}\frac{1}{4\pi}\mathcal{K}_1(z)\,|Y_\alpha|^2 \sum_\beta |Y_\beta|^2 \\ & \times\left[\frac{a_1^2+a_2^2}{\sqrt{c}} - \frac{2\sqrt{a_1a_2}(\sqrt{a_1c_1}+\sqrt{a_2c_2})}{(a_2-a_1)^2+(\sqrt{a_1c_1}+\sqrt{a_2c_2})^2}\left(a_1^{3/2}+a_2^{3/2}\right) + \mathcal{O}(Y'^2,\mu^2)\right].\end{aligned} \tag{3.6}$$

The first term in the parenthesis of eq. (3.7) arises from the pure N_1 and N_2 contributions, and yields inverse decay terms as in eq. (3.2). The second term is an interference effect, which is usually neglected in standard seesaw models. However, in the models we consider, this cannot be done, since both contributions are of the same order.

Factoring out the L-conserving contribution to the $\gamma^D_{N_{1,2,\alpha}}$ inverse decay terms, which up to $\mathcal{O}(Y'^2, \mu^2/M^2)$ terms is given by $\gamma^D_{N,\alpha} \equiv (\gamma^D_{N_1,\alpha} + \gamma^D_{N_2,\alpha})/2$, one finally obtains for the total washout

$$\begin{aligned}\gamma^{\mathrm{on}}_{\Delta L=2,\alpha} &= \frac{\gamma^D_{N,\alpha}}{4}\cdot 2\left(1 + 2\delta\sqrt{c} - \frac{1+3\delta\sqrt{c}}{1+\delta\sqrt{c}+\delta^2} + \mathcal{O}(\delta^2 c)\right)\left(1+\mathcal{O}(\delta\sqrt{c})\right) \\ &= \frac{\gamma^D_{N,\alpha}}{4}\cdot\left(2\,\delta^2 + \mathcal{O}(\delta^2 c)\right)\,,\end{aligned} \tag{3.7}$$

where the mass splitting enters in the parameter $\delta \equiv (M_2 - M_1)/\Gamma_1 = \mu/(M\sqrt{c}) \simeq (a_2 - a_1)/2\sqrt{c}$. There is an obvious cancellation in the parenthesis between the first two terms (pure N_1 and N_2 contributions) and the third one (N_1–N_2 interference), leaving a washout which vanishes in the L-conserving limit, $\delta \to 0$. Therefore, no matter how huge the "naive" inverse decay term in eq. (3.2) is, there is no washout in the L-conserving limit, as expected. Conversely, for a large mass splitting $\delta \gg 1$, the interference term becomes negligible, and one recovers the usual washout from inverse decays. This shows that the

[2]Here for simplicity we only consider the on-shell part, but it can be easily shown that the total $\Delta L = 2$ rate vanishes in the L-conserving limit, and therefore so does the off-shell part.

[3]Or equivalently, up to negligible terms, taking the residue of the integral in eq. (3.3) at the physical poles $\hat{a}_{1,2}$.

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interference term which is rightly neglected in typical seesaw models (see e.g. [27, 28]) can be essential in the context of approximate L-conserving frameworks.

Practically, since the washout vanishes in the L-conserving limit, it remains small as long as the L-violating perturbations, Y', $\mu_{1,2}$, remain small enough: it is suppressed as long as $\delta \ll 1$. The relevant washout parameter including the interference term turns out to be

$$K_\alpha^{\rm eff} \equiv K_\alpha \cdot \delta^2 \,. \tag{3.8}$$

For the numerical example above with $M \sim 1\,\text{TeV}$ and $K_\alpha \sim 10^{11}$, if $\delta \lesssim 10^{-4}$ one ends up with an effective washout $K_\alpha^{\rm eff} \simeq K_\alpha \cdot \delta^2 \lesssim 10^3$, which is suppressed enough to keep a substantial lepton asymmetry.

To sum up, in the class of models we are considering, the washout of the asymmetry can be suppressed by many orders of magnitude, granted that the degeneracy parameter δ is small enough. A splitting $\delta \ll 1$ is technically natural since it is controlled by L-violating parameters. As a matter of fact, the radiative and thermal corrections to δ are proportional to YY' and not to Y^2 as in typical seesaw scenarios. It is interesting to notice that, contrary to usual seesaw models with two RH neutrinos where the total washout is lower-bounded by the solar scale $K_1 + K_2 \gtrsim K_{\rm sol} \sim 9$, the effective washout in eq. (3.8) is lower-bounded only by terms quadratic in the small L-violating parameters Y' and μ/M. Therefore, $K_\alpha^{\rm eff} \ll K_{\rm sol}$ can be achieved in this case.

4 CP asymmetry

The *CP* asymmetry generated during the decays of N_i into the lepton flavour α is given by [20]

$$\epsilon_{i\alpha} = \frac{1}{8\pi} \sum_{j\neq i} \left\{ \frac{\text{Im}\left[h^*_{i\alpha} h_{j\alpha} \left(\sum_\gamma h^*_{i\gamma} h_{j\gamma}\right)\right]}{\sum_\beta |h_{i\beta}|^2} f_v^{i,j} + \frac{\text{Im}\left[h^*_{i\alpha} h_{j\alpha} \left(\sum_\gamma h_{i\gamma} h^*_{j\gamma}\right)\right]}{\sum_\beta |h_{i\beta}|^2} f_c^{i,j} \right\}, \tag{4.1}$$

where f_v is the usual L-violating self-energy and vertex loop factor and f_c is a L-conserving self-energy loop factor. In the limit we are interested in, namely $M_2 \simeq M_1$, only the self-energy correction is relevant, and we get $f_v^{2,1} \simeq -f_v^{1,2} \simeq f_c^{2,1} \simeq -f_c^{1,2} \equiv f_{\rm self}$ with [26, 29]

$$f_{\rm self} = \frac{a_2 - a_1}{(a_2 - a_1)^2 + \left(\sqrt{a_2 c_2} - \sqrt{a_1 c_1}\right)^2} \,. \tag{4.2}$$

It is instructive to express the *CP* asymmetry in terms of the parameters of the original Lagrangian, eqs. (2.1)–(2.2):

$$\epsilon_{1\alpha} = \epsilon_{2\alpha} \simeq -\frac{|Y_\alpha|^2}{4\pi} \left(\sin\alpha \frac{\mu_1 \mu_2}{2\mu M} + \frac{\sum_\beta \text{Im}(Y_\beta Y'^*_\beta e^{i\phi})}{\sum_{\beta'} |Y_{\beta'}|^2} \right) f_{\rm self} \,. \tag{4.3}$$

The above expression shows explicitly that the *CP* asymmetry in each flavour crucially depends on the L-violating parameters, and therefore vanishes when L is conserved.[4]

[4] In models with more fermion singlets, a term in the self-energy diagram may remain in the limit of L conservation, see e.g. $B - L$ conserving models of leptogenesis [30] or models with an extra N_3 [25, 31]. Unless an extra symmetry beyond L conservation is assumed to have a quasi-degeneracy between at least 3 fermion singlets, such a term turns out to be generically irrelevant for the generation of an asymmetry at the TeV scale due to lepton flavour equilibration [31].

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At the singularity $M_1 = M_2$, the self-energy correction above, eq. (4.2), has a regulator given by $\sqrt{a_2c_2} - \sqrt{a_1c_1}$. In usual seesaw scenarios, one typically has $|\sqrt{a_2c_2} - \sqrt{a_1c_1}| \simeq \max_i(\sqrt{a_ic_i})$ because the decay widths $\Gamma_{N_{1,2}}$ are very different from each other. However, in our case, while the $\sqrt{a_ic_i}$ do not vanish in the L-conserving limit, $\sqrt{a_2c_2} - \sqrt{a_1c_1}$ does. As a result $|\sqrt{a_2c_2} - \sqrt{a_1c_1}| \ll \min_i(\sqrt{a_ic_i})$. Therefore the resonance, which occurs for $a_2 - a_1 = \sqrt{a_2c_2} - \sqrt{a_1c_1}$, i.e. for $\delta = \delta_{\rm res} \simeq 2\,{\rm Re}(YY'^*e^{i\phi})/|Y|^2 - (\mu_2^2 - \mu_1^2)/2M\mu \ll 1$, is obtained for a mass splitting much smaller than the decay width. Around the resonance, where the *CP* asymmetry reaches its maximum value $\epsilon \sim 1$, the *CP* asymmetry is no longer suppressed by the small L-violating parameters. In practice in the following we will never need to use such a small mass splitting, since in this case, unless the *CP* phases are highly suppressed, the baryon asymmetry produced by leptogenesis would be far bigger than the observed value. Therefore, we will always stick to the natural range $|\sqrt{a_2c_2} - \sqrt{a_1c_1}| \ll a_2 - a_1 \ll \max_i(\sqrt{a_ic_i})$, or equivalently $\delta_{\rm res} \ll \delta \ll 1$. In this regime, the self-energy contribution is simply given by

$$f_{\rm self} \simeq \frac{1}{2\delta\sqrt{c}}. \tag{4.4}$$

It is remarkable that the conditions to get a large enough *CP* asymmetry and a suppressed washout are the same: $\delta \ll 1$.

Note that, using a different method, a regulator $\propto \sqrt{a_ic_i}$ was obtained in [18, 25], giving a resonance peak for $\delta \sim 1$. In the following we will stick to the regulator found in [26, 29], eq. (4.2), for two reasons. First, in [29], the self-energy contribution was calculated using two quantum field theory methods, and both converge to the regulator in eq. (4.2). Second, the very different oscillation formalism typically applying to the K^0–$\overline{K}^0$ and B^0–$\overline{B}^0$ systems was used in [20] to compute the leptogenesis self-energy diagram in the limit of small mass splitting and the same result was obtained. Note also that the results of [29], eq. (4.2), are obtained in the limit where the off-diagonal corrections to the propagator in the N_1–N_2 system are small compared to the diagonal ones. More precisely, in terms of the parameters of the model, the condition of ref. [29] reads $[(h^\dagger h)_{12} + (h^\dagger h)_{21}]/(32\pi^2) \ll (M_2 - M_1)/M_1$. In our framework, since the off-diagonal corrections are proportional to the L-violating parameters, i.e. $(h^\dagger h)_{12} + (h^\dagger h)_{21} \propto \mathcal{O}(Y'Y, \mu/M)$, this condition is clearly fulfilled.

5 Leptogenesis vs. rare lepton processes

It is well known that for RH neutrino masses of order the weak scale and Yukawa couplings of order few10^{-2}, rates for rare lepton processes such as $\mu \to e\gamma$ close to the present experimental bound, ${\rm Br}(\mu \to e\gamma) \lesssim 1.2 \times 10^{-11}$ [32], can be obtained [33, 34]. From the suppression of the washout in eq. (3.8) and the form of the *CP* asymmetry shown in eqs. (4.3) and (4.4), it is easy to convince oneself that successful leptogenesis is possible at the same time. Taking for example $M \sim 1\,{\rm TeV}$, $Y \sim 5 \times 10^{-2}$, $Y' \sim Y\mu/M \sim 10^{-10}$ (which may lead easily to the correct neutrino masses pattern), one gets $\delta \sim 10^{-5}$ and thus a *CP* asymmetry $\epsilon \sim$ few 10^{-6}. In this case, even though $K_\alpha \sim 10^{11}$, the effective washout in eq. (3.8) is reduced by a factor $1/\delta^2 \sim 10^{10}$, so that $K^{\rm eff} \sim 10^2$. Such values of the washout and of the *CP* asymmetry are in the right ballpark to explain the observed baryon asymmetry of the Universe.

To study this result more quantitatively it is convenient to use the Casas-Ibarra parametrization of the Yukawa couplings [35], which, in the 2 heavy neutrino case and for a normal hierarchy in the light sector, reads

$$h_{1\alpha} = \frac{\sqrt{M_1}}{v}\left(\sqrt{m_2}\,\cos z\, U^\star_{\alpha 2} + \sqrt{m_3}\,\sin z\, U^\star_{\alpha 3}\right)\,, \tag{5.1}$$

$$h_{2\alpha} = \frac{\sqrt{M_2}}{v}\left(-\sqrt{m_2}\,\sin z\, U^\star_{\alpha 2} + \sqrt{m_3}\,\cos z\, U^\star_{\alpha 3}\right)\,. \tag{5.2}$$

The inverted hierarchy is obtained simply by exchanging the indices $2 \to 1$ and $3 \to 2$. In this parametrization, $z = z_a + \mathrm{i} z_b$ is a complex angle, and U is the PMNS matrix. We will adopt the parametrization

$$U = \begin{pmatrix} c_{12}\,c_{13} & s_{12}\,c_{13} & s_{13}\,\mathrm{e}^{-\mathrm{i}\,\delta_{\rm CP}} \\ -s_{12}\,c_{23} - c_{12}\,s_{23}\,s_{13}\,\mathrm{e}^{\mathrm{i}\,\delta_{\rm CP}} & c_{12}\,c_{23} - s_{12}\,s_{23}\,s_{13}\,\mathrm{e}^{\mathrm{i}\,\delta_{\rm CP}} & s_{23}\,c_{13} \\ s_{12}\,s_{23} - c_{12}\,c_{23}\,s_{13}\,\mathrm{e}^{\mathrm{i}\,\delta_{\rm CP}} & -c_{12}\,s_{23} - s_{12}\,c_{23}\,s_{13}\,\mathrm{e}^{\mathrm{i}\,\delta_{\rm CP}} & c_{23}\,c_{13} \end{pmatrix} \cdot \mathrm{diag}(1, \mathrm{e}^{-\mathrm{i}\,\gamma/2}, 1), \tag{5.3}$$

where $s_{ij} \equiv \sin\theta_{ij}$, $c_{ij} \equiv \cos\theta_{ij}$. Neglecting experimental errors, we will use $\sin^2(2\theta_{12}) = 0.87$, $\sin^2(2\theta_{23}) = 0.92$ [32], $\sin^2(\theta_{13}) < 0.03$ (1σ) [36], whereas for the neutrino mass-squared differences we will take $\Delta m^2_{\rm sol} \simeq 7.59 \times 10^{-5}\,\mathrm{eV}^2$ and $\Delta m^2_{\rm atm} \simeq 2.43 \times 10^{-3}\,\mathrm{eV}^2$ [32]. Recall that with only two RH neutrinos the lightest active neutrino is massless. Therefore, if the neutrino mass hierarchy is normal, one has $m_1 = 0$, $m_2 = \sqrt{\Delta m^2_{\rm sol}}$ and $m_3 = \sqrt{\Delta m^2_{\rm sol} + \Delta m^2_{\rm atm}}$, whereas if it is inverted, $m_3 = 0$, $m_2 = \sqrt{\Delta m^2_{\rm atm}}$ and $m_1 = \sqrt{\Delta m^2_{\rm atm} - \Delta m^2_{\rm sol}}$.

In order to use the parametrization of eq. (5.1) in a regime corresponding to the model we consider, eqs. (2.1)–(2.2), let us invert eq. (2.7) to write the couplings Y and Y' as

$$Y_\alpha = \mathrm{e}^{-\mathrm{i}\phi/2}\sqrt{\frac{M}{2}\frac{\mathrm{e}^{z_b}}{v}}\,\mathrm{e}^{-\mathrm{i}\,z_a}\left(-\mathrm{i}\sqrt{m_2}\,U^*_{\alpha 2} + \sqrt{m_3}\,U^*_{\alpha 3}\right)\,, \tag{5.4}$$

$$Y'_\alpha = \mathrm{e}^{\mathrm{i}\phi/2}\sqrt{\frac{M}{2}\frac{\mathrm{e}^{-z_b}}{v}}\,\mathrm{e}^{\mathrm{i}\,z_a}\left(\mathrm{i}\sqrt{m_2}\,U^*_{\alpha 2} + \sqrt{m_3}\,U^*_{\alpha 3}\right) + \frac{Y_\alpha \mathrm{e}^{\mathrm{i}\phi}}{2}\left(\frac{\mu_2^2 - \mu_1^2}{2M\mu} + \mathrm{i}\lambda\right)\,. \tag{5.5}$$

It is clear that the parametrization above yields $Y' \ll Y$ if $\mathrm{e}^{-z_b} \ll 1$, the μ/M term being anyway suppressed by the approximate lepton number symmetry. In this limit, the *CP* asymmetry, eq. (4.3), takes the simple form:

$$\epsilon_\alpha = 4\,P_\alpha\,\frac{m_3 - m_2}{m_3 + m_2}\,\sin(2z_a)\,\frac{\mathrm{e}^{-2\,z_b}}{\delta} \tag{5.6}$$

where

$$P_\alpha = \frac{m_3|U_{\alpha 3}|^2 + 2\sqrt{m_3 m_2}\,\mathrm{Im}(U^*_{\alpha 2}U_{\alpha 3}) + m_2|U_{\alpha 2}|^2}{m_3 + m_2} \tag{5.7}$$

is the branching ratio of the decay of N_i into the flavour α, with $\sum_\alpha P_\alpha = 1$. *CP* violation in eq. (5.6) is controlled by the "leptogenesis phase" z_a and is maximal for $z_a = \pi/4$. It is interesting to notice that the *CP* asymmetry in the case of inverted hierarchy is suppressed compared to normal case by a factor $\sqrt{\Delta m^2_{\rm sol}/\Delta m^2_{\rm atm}}$. As for the $K_\alpha^{\rm eff}$ washout parameter, eq. (3.8), for $\delta \ll 1$ it is simply given by

$$K_\alpha^{\rm eff} \simeq \frac{M}{H(M)} \frac{|Y_\alpha|^2}{8\pi}\,\delta^2 = P_\alpha\, \frac{m_{3(2)} + m_{2(1)}}{2 m_\star}\; {\rm e}^{2\,z_b}\delta^2 \;\simeq \chi_1\, P_\alpha\, {\rm e}^{2\,z_b}\delta^2\,, \tag{5.8}$$

where $m_\star \simeq 1.08 \times 10^{-3}\,{\rm eV}$ and $\chi_1 \simeq 27\,(45)$ for normal (inverted) hierarchy. Had we neglected the interference term, we would have obtained a huge washout, $K_\alpha \simeq \chi_1\, e^{2\,z_b}$. Though the suppression of the washout $\propto \delta^2$ is crucial, it turns out that we will always remain in the regime of "strong" washout, defined as $K_\alpha^{\rm eff} \gtrsim 3$. In this regime, where the dependence on the initial conditions disappears, and where subtle effects like thermal corrections are irrelevant [27], we obtain finally for the baryon asymmetry

$$Y_B = \frac{12}{37} \sum_\alpha 0.5 \frac{\epsilon_\alpha}{(K_\alpha^{\rm eff})^{1.16}}\, Y_N^{\rm eq}(T \gg M) \simeq \chi_2 \left(\sum_\alpha P_\alpha^{-0.16} \right) \frac{e^{-4.32\, z_b}}{\delta^{3.32}} \sin{(2 z_a)}\,, \tag{5.9}$$

where $\chi_2 \simeq 8.7 \times 10^{-5}$ (5.4×10^{-7}) for normal (inverted) hierarchy. There is a mild dependence on the parameters of the PMNS matrix in $\sum_\alpha P_\alpha^{-0.16}$, which, barring large cancellations, varies between 3.7 and 4.8 for normal hierarchy and 3.6 and 4 for inverted. In the following we will constrain the baryon asymmetry to lie within the 3σ range of WMAP5, i.e. $8.1 \times 10^{-11} < Y_B < 9.5 \times 10^{-11}$ [37].

Let us now turn to the low-energy observables. The branching ratio for the rare lepton process $\mu \to e\gamma$ is given by [33]:

$$\begin{aligned} {\rm Br}(\mu \to e\gamma) &\simeq \frac{3\alpha}{32\pi}\left(2\, v^2\right)^2 \left| \frac{h_{1e} h_{1\mu}^\star}{M_1^2} + \frac{h_{2e} h_{2\mu}^\star}{M_2^2} \right|^2\,, \\ &\simeq \frac{3\alpha}{32\,\pi} \left(\frac{m_3 + m_2}{M} \right)^2 \; |g_{e\mu}|^2\, {\rm e}^{4 z_b}\,, \end{aligned} \tag{5.10}$$

where α is the fine structure constant, and

$$g_{e\mu} \equiv \frac{m_2\, U_{e2}^\star U_{\mu 2} + m_3\, U_{e3}^\star U_{\mu 3} + 2\,{\rm i}\sqrt{m_2 m_3}(U_{e3}^\star U_{\mu 2} - U_{e2}^\star U_{\mu 3})}{m_3 + m_2}\,. \tag{5.11}$$

As before, the replacement $3 \to 2$ and $2 \to 1$ should be made for the case of inverted hierarchy. The dependence on the parameters of the PMNS matrix is clearly larger for the branching ratio than for the baryon asymmetry. Here, $|g_{e\mu}|^2$ varies from 0.01 to 0.2 for the normal case, and from 0.06 to 0.7 in the inverted one, also granted that no large cancellations occur.

Using eq. (5.9) successful leptogenesis implies a relation between z_b and the degeneracy parameter δ, which we use to rewrite the branching ratio as

$$\begin{aligned} {\rm Br}(\mu \to e\gamma) &\simeq \frac{3\,\alpha}{32\,\pi} \left(\frac{m_3 + m_2}{M} \right)^2 \; |g_{e\mu}|^2 \left(\frac{\chi_2\, \left(\sum_\alpha P_\alpha^{-0.16}\right) \sin{(2 z_a)}}{Y_B^{\rm obs}\, \delta^{3.32}} \right)^{4/4.32} \\ &\simeq \chi_3 \times 10^{-12} \left(\frac{250\,{\rm GeV}}{M} \right)^2 \times \left(\frac{10^{-4}}{\delta} \right)^{3.1} \times \sin{(2 z_a)}^{0.93}\,, \end{aligned} \tag{5.12}$$

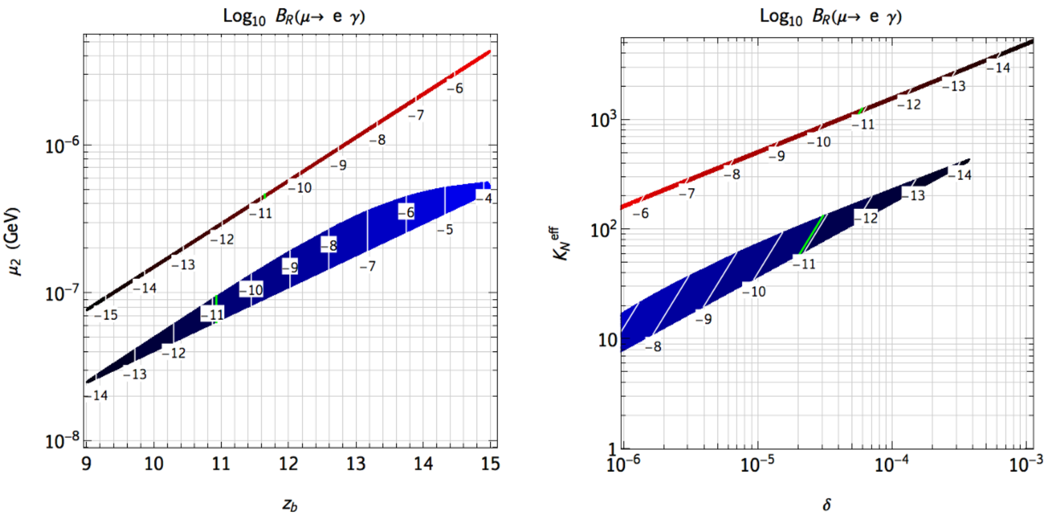

Figure 1. Contours of $\mathrm{Br}(\mu \to e\gamma)$ compatible with successful leptogenesis plotted against μ_2 and z_b (left panel), and $K^{\rm eff}$ and δ (right panel). The upper (lower) values correspond to the normal (inverted) hierarchical neutrino spectrum. The green line shows the present experimental limit $\mathrm{Br}(\mu \to e\gamma) \lesssim 1.2 \times 10^{-11}$.

with χ_3 taking values between 4.1×10^{-1} and 7.6 for the normal case and between 4.2×10^{-2} and 2.9×10^{-1} in the inverted one. From the above equations it is clear that the smaller the degeneracy parameter δ is, the bigger the Yukawa couplings can be without spoiling the success of leptogenesis, hence the larger the LFV rate $\mu \to e\gamma$ is. eq. (5.12) shows that for $M \sim 250\,\mathrm{GeV}$–$1\,\mathrm{TeV}$, observable $\mu \to e\gamma$ rates typically require $\delta \sim 10^{-3}$–10^{-5}. This corresponds to a mass splitting $(M_2 - M_1)/M_1 \sim 10^{-8}$–$10^{-10}$. Such a level of degeneracy is comparable to the one needed in usual resonant leptogenesis scenarios.

To check that our analytical result of eq. (5.12) is correct, it is necessary to integrate numerically the set of Boltzmann equations, as given in eq. (3.1). Figure 1 shows our numerical results for the simple case $\mu_1 = 0$, $M = 250\,\mathrm{GeV}$, $\delta_{\rm CP} = \gamma = 0$ and $\theta_{13} = 10^\circ$, taking for the leptogenesis phase z_a its maximum value $\pi/4$. In this simple case, as $\mu_1 = 0$, the phase α is unphysical and only the second term of eq. (4.3) contributes. The parameter μ_2 is directly responsible for the mass splitting between the RH neutrinos, while z_b sets the scale of the Yukawa couplings. Figure 1 displays, as a function of these 2 parameters, the iso-contours of the LFV rate compatible with successful leptogenesis for both normal and inverted neutrino mass hierarchies. A sphaleron freeze-out cut-off has been applied at the temperature $T \sim 130\,\mathrm{GeV}$ taking the Higgs boson mass $m_h = 120\,\mathrm{GeV}$ [39]. We find a good agreement with the estimate of eq. (5.12), up to the sphaleron freezeout which has not been taken into account in this equation. Also given in figure 1 are the LFV contour plot as a function of the parameters $K^{\rm eff}$ and δ. We observe that the branching ratio reached can saturate or even exceed the present experimental limit, $\mathrm{Br}(\mu \to e\gamma) \lesssim 1.2 \times 10^{-11}$, depicted as a green line in figure 1.

The $\mu \to e\gamma$ decay is not the only flavour violating process whose current experimental upper bound could be saturated in agreement with leptogenesis constraints. This is also

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possible for the τ radiative decays, which experimentally are constrained to the bounds $\mathrm{Br}(\tau \to \mu\gamma) < 4.5 \cdot 10^{-8}$ and $\mathrm{Br}(\tau \to e\gamma) < 1.1 \cdot 10^{-7}$ [32]. It turns out that the results of figure 1 can be applied to these rates for a degeneracy parameter δ about one order of magnitude smaller than for the $\mu \to e\gamma$ rate[5]. As explained in ref. [16] the observation of several radiative decays would allow to overconstrain this model. In the type-I seesaw models, other channels like $\mu \to eee$ or $\tau \to 3l$ are expected to be more suppressed than the radiative decays, see e.g. ref. [34].

6 Softly broken L case: $Y'_\alpha = 0$

At the end of the previous section, by setting μ_1 to 0 in figure 1, we numerically considered a simple case where the *CP* asymmetry is proportional to the Y' couplings. It is interesting to discuss what happens if we take $Y' = 0$ instead. This is the situation of the usual inverse seesaw models [9]–[12] where L is assumed to be softly broken, and where in full generality n pairs of $N_{1,2}$ are usually assumed. In such a framework the Yukawa coupling matrix Y in eqs. (2.1)–(2.3) is an $n \times 3$ matrix, $\mu_{1,2}$ are $n \times n$ matrices and Y' is a null $n \times 3$ matrix. It must be stressed that in this case too leptogenesis can be compatible with observable flavour violation rates. This can be seen from the $n = 1$ case above which, through the first term of eq. (4.3), gives the *CP* asymmetry

$$\epsilon_{1\alpha} = \epsilon_{2\alpha} \simeq -\frac{|Y_\alpha|^2}{4\pi} \sin\alpha \, \frac{\mu_1\mu_2}{2\mu M} f_{\mathrm{self}} \overset{\delta \ll 1}{\simeq} -\frac{|Y_\alpha|^2}{16\pi} \frac{\mu_1\mu_2}{\mu^2} \sin\alpha \,, \tag{6.1}$$

where eq. (4.4) was used in the last equality. It is worth mentioning again that the *CP* asymmetry is not suppressed by the small L-violating entries in the limit of small splitting $\delta \ll 1$. The larger are the Y_α couplings, the larger is the $\mu \to e\gamma$ rate and the larger is the *CP* asymmetry. Note that a non-vanishing value of both μ_1 and μ_2 is required in order to have a non-zero *CP* asymmetry, since otherwise the phase α can be rotated away. As for the washout, it is again suppressed as soon as $\delta \ll 1$, eq. (3.7).

The generalization to the inverse seesaw case with $n > 1$ pairs of $N_{1,2}$ is straightforward. In this case the N_1's and N_2's form n pairs of quasi-degenerate states and, as in eqs. (3.7) and (6.1), each pair gives an unsuppressed *CP* asymmetry and a small washout as soon as the mass splitting is smaller than the decay width. Note that when $Y'_\alpha = 0$, $n \geq 2$ is required in order to give at least 2 neutrino masses. The detailed study of this case is beyond the scope of our paper.

7 Conclusion

We have shown that, in the context of approximately L-conserving seesaw models, it is rather easy to generate large lepton flavour effects in agreement with successful leptogenesis. These models provide all the necessary ingredients. They involve large Yukawa couplings without leading to large neutrino masses, as required to induce large rates. They predict a small right-handed neutrino mass splitting protected from large radiative corrections by

[5]The Yukawa couplings in the μ-e channel should be somewhat suppressed with respect to the ones of the other channel(s) for the more stringent $\mu \to e\gamma$ constraint to be satisfied.

the approximate L-symmetry. If the splitting is sufficiently small, $\delta \ll 1$, the washout is suppressed, proportionally to the small L-violating entries. And, finally, from the same condition, $\delta \ll 1$, the suppression of the CP asymmetry by the small L-violating entries is compensated by its resonant behaviour. In this way the experimental upper bound of any of the rare lepton processes $\mu \to e\gamma$, $\tau \to \mu\gamma$, $\tau \to e\gamma$ can be saturated in agreement with successful leptogenesis. We emphasize that this pattern can work with only 2 RH neutrinos and that moreover it does not require to assume any flavour symmetry. Only a lepton number assignment for the various RH neutrinos has to be assumed for lepton number to be approximately conserved.

To illustrate these results we have considered a minimal 2 RH neutrino model which has the additional virtue that from the knowledge of the neutrino mass matrix one can predict the rates of any rare lepton process up to an overall normalization scale. Therefore this model can be overconstrained. To our knowledge, if one does not assume any extra particles beside the RH neutrinos, there is no neutrino mass and leptogenesis seesaw model more testable than this one.

The observation of rare lepton processes is not the only way to probe the inverse seesaw framework. Firstly, a deviation from unitarity in the light neutrino mixing matrix testable in a future neutrino factory can arise [40, 41]. Secondly, the unsuppressed active-sterile mixing angles allow in principle for a non-negligible RH neutrino production cross-section at the LHC in the channel $pp \to W^* \to N\ell$, with the RH neutrino subsequently decaying to lepton and W. However, the almost background-free signal in two like-sign leptons is suppressed by the smallness of neutrino masses [13, 42]. On the other hand, the recently noticed trilepton channel plus missing energy might be observable [42]. Concerning lepton-flavour-violating channels, the current bound from $\mu \to e\gamma$ seems to seriously limit the reach at the LHC [43].

Our results also apply to the usual inverse seesaw models where L is softly broken, with n pairs of fermion singlets and $Y' = 0$. Similarly it also applies to the 3 RH neutrino case where, if L conservation is approximate, two RH neutrinos are naturally quasi-degenerate. In particular they apply to the 3 RH neutrino models considered in refs. [13, 14, 16]. For the third RH neutrino, one has just to make sure that its mass is either sizably heavier than the two others or that at least one of its Yukawa coupling is suppressed enough.

Finally it is worth mentioning that our main result about the washout suppression also applies to the Type-III seesaw models with the same Yukawa coupling matrix and heavy state mass matrix. However, leptogenesis suffers in this case from the high degree of thermalization of the fermion triplets by the gauge interactions [44]. It would be interesting to study more extensively this case to determine which region of the parameter space is allowed.

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