Minimal Resonant Leptogenesis and Lepton Flavour Violation

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Abstract. We discuss minimal non-supersymmetric models of resonant leptogenesis, based on an approximate flavour symmetries. As an illustrative example, we consider a resonant $\tau$-leptogenesis model, compatible with universal right-handed neutrino masses at the GUT scale, where the required heavy-neutrino mass splittings are generated radiatively. In particular, we explicitly demonstrate, how a minimum number of three heavy Majorana neutrinos is needed, in order to obtain successful leptogenesis and experimentally testable rates for processes of lepton flavour violation, such as $\mu \to e\gamma$ and $\mu \to e$ conversion in nuclei.

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INTRODUCTION

The observed baryon asymmetry in the Universe (BAU), which amounts to a baryon-to-photon ratio of number densities $\eta_B \approx 6.2 \times 10^{-10}$ [1, 2], provides one of the strongest pieces of evidence for physics beyond the Standard Model (SM) [3]. One interesting scenario for explaining the BAU is leptogenesis [4]. A potentially interesting alternative to GUT-scale leptogenesis is the framework of low-scale resonant leptogenesis (RL) [5, 6]. Within this framework, the lowering of the scale may rely on a dynamical mechanism, in which heavy-neutrino self-energy effects [7] on the leptonic asymmetry become dominant [8] and get resonantly enhanced [5], when a pair of heavy Majorana neutrinos has a mass difference comparable to the decay widths.

Flavour effects due to heavy-neutrino Yukawa couplings [9] play an important role in models of RL [10, 11]. In particular, in [10] a scenario was put forward, called resonant $\tau$-leptogenesis (R$\tau$L), in which the BAU originates from a $\tau$-lepton asymmetry, resonantly produced by quasi-in-equilibrium decays of heavy Majorana neutrinos. In a R$\tau$L model, the generated excess in the $L_\tau$ number will be converted into the observed BAU, provided the $L_\tau$-violating reactions are not strong enough to wash out such an excess. In such a scenario, the heavy Majorana neutrinos can be as light as 100 GeV and have sizeable couplings to two of the charged leptons, specifically to electrons and muons. Consequently, depending on the flavour dynamics of heavy neutrino Yukawa coupling effects, phenomenologically testable models of RL can be built that could be probed at the LHC or in low energy experiments of lepton number violation (LNV) and lepton flavour violation (LFV), such as $\mu \to e\gamma$ and neutrinoless double beta decay. As we demonstrated in [12], R$\tau$L models can be constructed with a degenerate heavy neutrino mass spectrum at the GUT scale. In a minimal non-supersymmetric framework, the splittings required for successful RL are generated radiatively [13] by renormalization group (RG) evolution and can therefore be naturally comparable to the decay widths of the heavy neutrinos.

$\tau$-LEPTON FLAVOUR SYMMETRY

The leptonic Yukawa and Majorana sectors of the SM symmetrically extended with one singlet right-handed neutrino $v_{iR}$ per $i$ family (with $i = 1, 2, 3 = e, \mu, \tau$) are given by the Lagrangian

$$-\mathcal{L}_{Y,M} = \bar{L}_i \Phi h^t l_R + \bar{L}_i \Phi h^t v_R + \bar{v}^C_R m_M v_R + H.c.,$$

where $\Phi$ is the SM Higgs doublet and $\bar{\Phi} = i\tau_2 \Phi^*$ its isospin conjugate. Moreover, we have suppressed the generation index $i$ from the left-handed doublets $L_i = (v_{iL}, l_{iL})^T$, the right-handed charged leptons $l_{iR}$ and the right-handed neutrinos $v_{iR}$, while ordinary multiplication between vectors and matrices is implied.
To obtain a phenomenologically relevant model in this minimal setup, at least 3 singlet heavy Majorana neutrinos $v_{1,2,3 R}$ are needed and these have to be nearly degenerate in mass [10]. To ensure the latter, we assume that to leading order, the singlet Majorana sector is SO(3) symmetric, i.e.

$$m_M = m_N I_3 + \Delta m_M,$$

where $I_3$ is the $3 \times 3$ identity matrix and $\Delta m_M$ is a general SO(3)-breaking matrix induced by RG effects. As we will discuss below, compatibility with the observed light neutrino masses and mixings requires that $(\Delta m_M)_{ij}/m_N \lesssim 10^{-7}$, for electroweak-mass Majorana neutrinos, i.e. for $m_N \approx 0.1$–1 TeV. We will explicitly demonstrate, how such an SO(3)-breaking matrix $\Delta m_M$, of the required order, can be generated radiatively via the RG evolution of the right-handed neutrino mass matrix $m_M$ from the GUT scale $M_X \approx 2 \times 10^{16}$ GeV to the mass scale of the right-handed neutrinos $m_N$.

In the physical charged-lepton mass basis, the SO(3) symmetry imposed on the singlet Majorana sector at the GUT scale $M_X$ gets explicitly broken by a set of neutrino Yukawa couplings to the subgroup of lepton symmetries: $U(1)_{L_e+L_{\mu}} \times U(1)_{L_\tau}$. The flavour charge assignments that give rise to such a breaking are presented in Table 1.

As a consequence of the $U(1)_{L_e+L_{\mu}} \times U(1)_{L_\tau}$ symmetry, the neutrino Yukawa coupling matrix takes on the general form:

$$h^V = \begin{pmatrix} 0 & ae^{-i\pi/4} & ae^{i\pi/4} \\ 0 & be^{-i\pi/4} & be^{i\pi/4} \\ 0 & 0 & 0 \end{pmatrix} + \begin{pmatrix} \epsilon_e & 0 & 0 \\ \epsilon_\mu & 0 & 0 \\ \epsilon_\tau & \kappa_1 e^{i(\pi/4-\gamma_1)} & \kappa_2 e^{i(\pi/4-\gamma_2)} \end{pmatrix},$$

where the second term vanishes, if the symmetry is exact. In this symmetric limit, the light neutrinos remain massless to all orders in perturbation theory, whilst $a$ and $b$ are free unconstrained complex parameters. The phases accompanying these parameters in (3) are simply chosen for convenience to maximize the lepton asymmetry in leptogenesis when $a$ and $b$ are real. In order to give masses to the light neutrinos, the second term describes the departure from the flavour symmetric limit with $|\epsilon_e, \epsilon_\mu, \epsilon_\tau|, \kappa_1 = \kappa_2 \ll |a|, |b|$ and the phases $\gamma_1, \gamma_2$ are unrestricted.

It is important to notice that the flavour structure of the neutrino Yukawa couplings $h^V$ is preserved through the RG evolution, as long as the flavour symmetry is only weakly broken. In detail, RG effects violate the SO(3)-invariant form of the Majorana mass matrix $m_M(M_X) = m_N I_3$ by the 3-by-3 matrix $\Delta m_M$. In the leading-log approximation, the SO(3)-breaking matrix $\Delta m_M$ receives the following radiative corrections:

$$\Delta m_M = -\frac{m_N}{8\pi^2} \ln \left( \frac{M_X}{m_N} \right) \text{Re} \left[ h^V(M_X) h^V(M_X) \right].$$

Given the form invariance of $h^V$ under RG effects, we may therefore define all input parameters at the right-handed neutrino mass scale $m_N$, where the matching with the light neutrino data is performed.

We may now determine the Yukawa parameters $(a,b,\epsilon_e,\epsilon_\mu,\epsilon_\tau)$, in terms of the light-neutrino mass matrix $m^V$ in the positive and diagonal charged lepton Yukawa basis. To do so, we first notice that the chosen symmetry $U(1)_{L_e+L_{\mu}} \times U(1)_{L_\tau}$ ensures the vanishing of the light neutrino mass matrix $m^V$. However, the symmetry-breaking parameters induce a non-zero light neutrino mass matrix $m^V$, which to leading order in these parameters is given by [11, 12]

$$m^V = -\frac{\nu^2}{2} h^V m_M h^V h^T = \frac{\nu^2}{2m_N} \left( h^V \Delta m_M h^T \right) \left( \frac{m_N}{m_N} \right) - h^V h^T,$$

$$= -\frac{\nu^2}{2m_N} \left( \frac{\Delta m_N}{m_N} a^2 - \epsilon_e^2 \Delta m_N ab - \epsilon_e \epsilon_\mu \Delta m_N b^2 - \epsilon_\mu \epsilon_\tau \Delta m_N \right),$$

$$= -\frac{\nu^2}{2m_N} \left( \frac{\Delta m_N}{m_N} ab - \epsilon_e \epsilon_\mu \right) \left( \frac{\Delta m_N}{m_N} b^2 - \epsilon_\mu \epsilon_\tau \right),$$

$$= -\frac{\nu^2}{2m_N} \left( \Delta m_N ab - \epsilon_e \epsilon_\mu \right),$$

TABLE 1. Flavour charge assignments for the breaking $U(1)_{L_e+L_{\mu}} \times U(1)_{L_\tau}$.

| $L_e, e_R$ | $L_\mu, \mu_R$ | $L_\tau, \tau_R$ | $\nu_1$ | $\nu_2 \pm i\nu_3$ |
|----------|-------------|-------------|------|-----------------|
| $U(1)_{L_e+L_{\mu}}$ | $+1$ | $+1$ | $0$ | $0$ | $\pm 1$ |
| $U(1)_{L_\tau}$ | $0$ | $0$ | $+1$ | $0$ | $0$ |
where \( v = 2M_W/g_w = 245 \text{ GeV} \) is the vacuum expectation value of the SM Higgs field \( \Phi \). In deriving the last equation in (5), we have also assumed that \( \sqrt{\frac{m_N}{m_N}} \kappa_{1,2} \ll \varepsilon_{e,\mu,\tau} \), where \( \Delta m_N \) stands for the expression

\[
\Delta m_N \equiv 2(\Delta m_{12})_{23} + i \left[ (\Delta m_{12})_{33} - (\Delta m_{12})_{22} \right] = -\frac{m_N}{8\pi^2} \ln \left( \frac{M_\chi}{m_N} \right) \left[ 2\kappa_1 \kappa_2 \sin(\gamma_1 + \gamma_2) + i(\kappa_2^2 - \kappa_1^2) \right].
\]

As a consequence of the flavour symmetry \( U(1)_{L^e + L^\mu} \times U(1)_{L^\tau} \), the symmetry-violating parameters \( \varepsilon_{e,\mu,\tau} \) and \( \kappa_{1,2} \) enter the tree-level light neutrino mass matrix \( m^v \) in (5) quadratically. This in turn implies that for electroweak-scale heavy neutrinos \( m_N \sim v \), the symmetry-breaking Yukawa couplings need not be much smaller than the electron Yukawa coupling \( h_e \sim 10^{-6} \). Moreover, one should observe that only a particular combination of SO(3)-violating terms \( (\Delta m_{ij})_{ij} \) appears in \( m^v \) through \( \Delta m_N \). Nevertheless, for electroweak-scale heavy neutrinos with mass differences \( |\Delta m_N|/m_N \lesssim 10^{-7} \), one should have \( |a|, |b| \lesssim 10^{-2} \) to avoid getting too large light neutrino masses much above 0.5 eV.

Finally, we note that one-loop Z- and Higgs-boson effects can induce additional contributions to the light neutrino mass matrix \( m^v \) [14]. As well as being loop suppressed, these radiative contributions are proportional to the symmetry-breaking parameters and can therefore be neglected for the R \( \tau \) L scenario under study [11].

Given the analytic form (5) of the light neutrino mass matrix, we may directly compute the neutrino Yukawa couplings \( (a, b, \varepsilon_e, \varepsilon_\mu, \varepsilon_\tau) \), as functions of \( m_N \), the phenomenologically constrained neutrino mass matrix \( m^v \) and the symmetry-breaking parameters \( \kappa_{1,2} \) and \( \gamma_{1,2} \):

\[
a^2 = \frac{2m_N}{v^2} \frac{8\pi^2}{\ln(M_\chi/m_N)} \left( \frac{m^v_{11} - (m^v_{13})^2}{m^v_{33}} \right) \left[ 2\kappa_1 \kappa_2 \sin(\gamma_1 + \gamma_2) + i(\kappa_2^2 - \kappa_1^2) \right]^{-1},
\]

\[
b^2 = \frac{2m_N}{v^2} \frac{8\pi^2}{\ln(M_\chi/m_N)} \left( \frac{m^v_{22} - (m^v_{23})^2}{m^v_{33}} \right) \left[ 2\kappa_1 \kappa_2 \sin(\gamma_1 + \gamma_2) + i(\kappa_2^2 - \kappa_1^2) \right]^{-1},
\]

\[
\varepsilon_e^2 = \frac{2m_N}{v^2} \frac{m^v_{13}^2}{m^v_{33}}, \quad \varepsilon_\mu^2 = \frac{2m_N}{v^2} \frac{m^v_{23}^2}{m^v_{33}}, \quad \varepsilon_\tau^2 = \frac{2m_N}{v^2} \frac{m^v_{33}}{m^v_{33}}.
\]

Since the approximate light-neutrino mass matrix \( m^v \) in (5) has rank 2, the lightest neutrino mass eigenstate \( \nu_1 \) will be massless in this approximation. The relations given in (7) will be used to obtain numerical estimates of the BAU, in terms of \( m_N \) and the symmetry-breaking parameters \( \kappa_{1,2} \) and \( \gamma_{1,2} \), for both normal and inverted hierarchy scenarios of light neutrinos.

**RESONANT LEPTOGENESIS**

Within the framework of leptogenesis, a net non-zero leptonic asymmetry results from the CP-violating decays of the heavy Majorana neutrinos \( N_\alpha \) into the left-handed charged leptons \( l_L^\pm \) and light neutrinos \( \nu_L \). Consequently, we have to calculate the partial decay width of the heavy Majorana neutrino \( N_\alpha \) into a particular lepton flavour \( l \),

\[
\Gamma_{\alpha l} = \Gamma(N_\alpha \rightarrow l_L^+ + W^+) + \Gamma(N_\alpha \rightarrow \nu_L + Z, H).
\]

In RL models, resumming the absorptive parts of the heavy Majorana-neutrino self-energy transitions \( N_\beta \rightarrow N_\alpha \) plays an important role in the computation of \( \Gamma_{\alpha l} \) [5, 6].

The leptonic asymmetries for each individual lepton flavour can be expressed in terms of the resummed neutrino Yukawa couplings \( h^v_{\alpha l} [6, 15] \) as

\[
\delta_{\alpha l} \equiv \frac{\Gamma_{\alpha l} - \Gamma_{\alpha l}^C}{\sum_{l=e,\mu,\tau} \left( \Gamma_{\alpha l} + \Gamma_{\alpha l}^C \right)} = \frac{\left| h^v_{\alpha l} \right|^2 - \left| h_{\alpha l}^C \right|^2}{\left( h^v_{\alpha l} h_{\alpha l}^v \right)_{\alpha\alpha} + \left( h_{\alpha l}^C h_{\alpha l}^C \right)_{\alpha\alpha}}.
\]

The analytic results for the leptonic asymmetries \( \delta_{\alpha l} \) simplify considerably in the 2-heavy neutrino mixing limit [5, 6],

\[
\delta_{\alpha l} \approx \frac{\text{Im} \left( (h^v_{\alpha l} h_{\beta l}) (h^v_{\beta l} h^v_{\alpha l})_{\beta\beta} \right)}{(h^v_{\alpha l} h_{\alpha l}^v)_{\alpha\alpha} (h^v_{\beta l} h_{\beta l})_{\beta\beta}} \frac{(m_{N_\alpha} - m_{N_\beta}) m_{N_\alpha} \Gamma_{\alpha l}^{(0)}}{(m_{N_\alpha} - m_{N_\beta})^2 + m_{N_\alpha}^2 \Gamma_{\alpha l}^{(0)}},
\]
of the leptonic asymmetries. A detailed comparison of the predictions obtained by different approaches to resonant
technology is given in [12].

The Boltzmann equations for the heavy Majorana neutrino and lepton-number densities, \( \eta_{N,1,2,3}, \eta_{e,\mu,\tau} \) normalized to the photon number density \( n_\gamma \) are given by [11, 12]

\[
\frac{d\eta_{N_i}}{dz} = -\frac{z}{n_\gamma H_N} \left( \frac{\eta_{N_i}}{\eta_{N_i}^{eq}} - 1 \right) \gamma_{N_i}^L, \\
\frac{d\eta_{L_{\alpha}}}{dz} = \frac{z}{n_\gamma H_N} \left[ \sum_{\alpha=1}^{3} \left( \frac{\eta_{N_\alpha}}{\eta_{N_\alpha}^{eq}} - 1 \right) \delta_{\alpha L_{\alpha}} + \frac{3}{2} \eta_{\ell_{\alpha}} \sum_{k=e,\mu,\tau} \left( \gamma_{L_k \Phi}^{j_{\alpha,\Phi}} + \gamma_{L_k \Phi}^{j_{\alpha,\Phi}^*} \right) - \frac{2}{3} \sum_{k=e,\mu,\tau} \eta_{\ell_k} \left( \gamma_{L_k \Phi}^{j_{\ell,\Phi}} - \gamma_{L_k \Phi}^{j_{\ell_k,\Phi}^*} \right) \right],
\]

where \( \alpha = 1, 2, 3 \) and \( L = e, \mu, \tau \). In addition, \( H_N \approx 17 \times m_N^2/M_p \) is the Hubble parameter at \( T = m_N \), where \( M_p = 1.2 \times 10^{19} \) GeV is the Planck mass. The \( T \)-dependence of the BEs (13) and (14) is expressed in terms of the dimensionless parameter \( z = H_N \). In the equilibrium number density of the heavy Majorana neutrino \( N_\alpha \), normalized to the number density of photons. The BEs (13) and (14) include the collision terms \( \gamma(X \to Y) \equiv \gamma_{L_{\alpha} \Phi}^{j_{\alpha,\Phi}} \) for the decays \( N_\alpha \to L_{\alpha} \Phi \), as well as the \( \Delta L = 0.2 \) resonant scattering processes \( L_{\alpha} \Phi \to L_{\alpha} \Phi \) and \( L_{\alpha} \Phi \to L_{\alpha} \Phi^* \) [11].

The lepton asymmetries are then partially converted via \((B+L)\)-violating sphaleron processes [16]. For temperatures \( T \) larger than the critical temperature \( T_c \approx 135 \) GeV of the electroweak phase transition, the conversion of the total lepton-to-photon ratio \( \sum_{L=e,\mu,\tau} \eta_{L_{\alpha}} \) to the baryon-to-photon ratio \( \eta_{\bar{B}} \) at \( T_c \) is given by \( \eta_{\bar{B}} = \frac{28}{3} \sum_{L=e,\mu,\tau} \eta_{L_{\alpha}} \).
Heavy Majorana-neutrino loop effects may induce sizeable LFV couplings to the photon and the Z boson. These couplings give rise to LFV decays, such as $\mu \rightarrow e\gamma$ [17], $\mu \rightarrow eee$ [18] and $\mu \rightarrow e$ conversion in nuclei. In the RRL model, only two of the right-handed neutrinos have appreciable $e$- and $\mu$-Yukawa couplings $a, b \approx 10^{-2}$ for $\kappa_{1,2} \approx 10^{-5} - 10^{-3}$. The LFV branching ratio $B(\mu \rightarrow e\gamma)$ can be expressed as [12]

$$B(\mu \rightarrow e\gamma) \approx 8.0 \cdot 10^{-4} \times g\left(\frac{m_N}{m_W}\right) \frac{v^4}{m_W^2} a^2 b^2,$$

with the loop function $g(x)$, possessing the limits $g \rightarrow 1$ for $m_N \gg m_W$ and $g \rightarrow 1/16$ for $m_N = m_W$. This theoretical prediction can be favourably contrasted with the current experimental upper limit $B(\mu \rightarrow e\gamma) < 2.4 \cdot 10^{-12}$ [19].

After dilution of this baryon asymmetry $\eta_B$ by standard photon interactions until the recombination epoch, $\eta_B \approx 1/27 \eta_B^0$, the theoretical prediction can now be compared with the current observational value for the baryon-to-photon asymmetry [2]: $\eta_B^{\text{obs}} = (6.20 \pm 0.15) \times 10^{-10}$.

In the limit of strong washout, the baryon asymmetry $\eta_B$ becomes relatively independent of the initial values of $\eta_B^0$ and $\eta_B^0$. In this case, the Boltzmann equations can be solved analytically and the resulting BAU can be estimated to be [12]

$$\eta_B \approx \frac{-3 \cdot 10^{-2} \delta_\tau}{K_\tau \min\left(m_N/T_\tau, 1.25 \ln(25K_\tau)\right)},$$

with the total $\tau$ asymmetry $\delta_\tau = \sum a_\alpha \delta \alpha \tau$ and the total $\tau$ asymmetry washout factor $K_\tau = \sum a_\alpha \Gamma_{N_\alpha \rightarrow L_\tau \phi}/H_N$. This formula provides a good approximation of the generated BAU in the RRL scenario within the regime $K_\tau \gtrsim 5$ for a right-handed neutrino mass scale of the order of the EW scale. Hence, to account for the observed BAU, a $\tau$ lepton asymmetry $\delta_\tau \gtrsim 10^{-7}$ is required.

**LEPTON FLAVOUR VIOLATION AND $\theta_{13}$**

FIGURE 2. Baryon asymmetry $\eta_B$ (blue contours) and the LFV observables $B(\mu \rightarrow e\gamma)$ and $B_{\mu e}(13,15,17)$ (red contours) as functions of $\kappa_1$ and $\kappa_2$ in the RRL model with $m_N = 120$ GeV, assuming a normal (left) and inverse (right) light neutrino mass spectrum. The remaining parameters are: $\gamma_1 = 3\pi/8$, $\gamma_2 = \pi/2$, $\phi_1 = \pi$, $\phi_2 = 0$, $\text{Re}(a) > 0$. The neutrino oscillation parameters are set at their best fit values, with $\sin^2 \theta_{13} = 0.027$. The blue shaded regions denote the parameter space where the baryon asymmetry is larger than the observed value. The green shaded areas labeled as ‘$\Delta m^2 = 0.25$’ and ‘$0.10$’ indicate the parameter space where the inversion of the light-neutrino mass matrix is violated at the 25% and 10% level, respectively.
Minimal low-scale seesaw scenarios of resonant leptogenesis are a class of models with interesting implications for observables of charged LFV, such as $\mu \rightarrow e\gamma$ and $\mu \rightarrow e$ conversion in nuclei. Here we have discussed aspects of
the \( R_{\tau L} \) realization of resonant leptogenesis, where the observed BAU originates from an individual \( \tau \) lepton-number asymmetry which gets resonantly enhanced via the out-of-equilibrium decays of nearly degenerate heavy Majorana neutrinos. The required mass splittings are generated radiatively via RG effects with universal right-handed neutrino masses at the GUT scale.

Analogous to the \( R_{\tau L} \) case discussed here, \( R_{e L} \) and \( R_{\mu L} \) scenarios with approximately protected \( e, \mu \) asymmetries, respectively, may be constructed [12]. Heavy Majorana neutrinos in such \( R_{\ell L} \) scenarios can be as light as 100 GeV, whilst their couplings to two of the charged leptons may be large, so as to lead to LFV effects that could be tested by the MEG and the COMET/PRISM experiments. In the \( R_{\ell L} \) model with a normal light neutrino mass hierarchy, there is a sizeable model parameter space with successful leptogenesis and large LFV process rates, with \( B(\mu \rightarrow e\gamma) \approx 10^{-12} \). This prediction is largely independent of \( \sin^2 \theta_{13} \) and the other light neutrino oscillation parameters. On the other hand, in the \( R_{\tau L} \) model with inversely hierarchical light neutrinos, \( B(\mu \rightarrow e\gamma) \) is linearly proportional to \( \sin^2 \theta_{13} \), and can be enhanced by more than one order of magnitude compared to the normal hierarchy case for \( \sin^2 \theta_{13} \) close to its upper experimental 2\( \sigma \) limit. Unfortunately, the generated baryon asymmetry \( \eta_B \) is suppressed in this scenario, and to test the viable parameter space for successful leptogenesis would require an experiment for \( \mu \rightarrow e \) conversion in nuclei which is sensitive to \( B_{\mu e} \approx 10^{-17}–10^{-16} \).

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