LEPTON FLAVOUR EFFECTS AND RESONANT LEPTOGENESIS

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Minimal models exist in which thermal, resonant leptogenesis occurs at the electroweak scale. The effects of individual lepton flavours play a crucial role and allow successful leptogenesis with large couplings between some of the charged leptons and the heavy Majorana neutrinos. This leads to potentially observable signals for low energy experiments (such as MEG, MECO and PRIME) and future linear colliders.

1 Introduction

If the mass difference between two or more heavy Majorana neutrinos is much smaller than their masses then the CP-asymmetry in their decays occurs primarily through self-energy effects ($\varepsilon$-type CP-violation) rather than vertex effects ($\varepsilon'$-type CP-violation). If this mass difference is of the order of the heavy neutrino decay widths then the CP-asymmetry can become resonantly enhanced and may be of order 1. This resonant enhancement can be exploited in models where the mass scale of the heavy Majorana neutrinos is as low as the electroweak scale (i.e. $\sim 250$ GeV) allowing successful thermal leptogenesis in complete accordance with current light neutrino data. Such models avoid the severe constraints on the scale of the heavy Majorana neutrinos arising when they have a hierarchical mass spectrum and thermal leptogenesis is required; these constraints arise because the lightest heavy neutrino must have a mass lower than the reheat temperature ($T_{RH}$) where $T_{RH} \lesssim 10^9$ GeV. A stringent limit on the reheat temperature is necessary in models incorporating supergravity to avoid the overproduction of gravitinos in the early Universe, which would disrupt the nucleosynthesis of the light elements. In addition, the class of models under study should be testable through experiments sensitive to lepton flavour and number violation and may lead to the production of heavy Majorana neutrinos at a high energy linear collider.

We will work within the framework of the Standard Model (SM) supplemented by 3 generations of gauge-singlet right-handed neutrinos $\nu_R$. As in the usual see-saw scenario, these gauge singlets are permitted to have a Majorana mass, $M_S$, violating lepton number by two units. In the basis where the charged lepton Yukawa couplings are diagonal, the relevant part of the Lagrangian reads

$$-\mathcal{L}_{M,Y} = \sum_{i,j=1}^{3} \frac{1}{2} (\bar{\nu}_{iR})^C (M_S)_{ij} \nu_{jR} + \sum_{k=\epsilon,\mu,\tau} \left( \hat{h}^{l}_{kk} \bar{L}_k \Phi l_R + \sum_{l=1}^{3} h^{\nu R}_{kl} \bar{L}_k \tilde{\Phi} \nu_{lR} \right) + \text{H.c.} , \quad (1)$$

where $\tilde{\Phi} = i \tau_2 \Phi^*$ is the isospin conjugate of the Higgs doublet $\Phi$ and $\tau_2$ is the usual Pauli matrix. $L$ and $l_R$ represent lepton SU(2)$_L$ doublets and singlets respectively, $h^{\nu R}$ is the matrix of neutrino Yukawa couplings and $h^l$ is the matrix of charged lepton Yukawa couplings.

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In the context of leptogenesis, the Sakharov requirement for a departure from thermal equilibrium translates into the requirement that the heavy neutrino decay rate is smaller than the expansion rate of the Universe \( H(T) \). This can be quantified by the parameter \( K_{N_i}^l \), where

\[
K_{N_i}^l \equiv \left[ \Gamma(N_i \rightarrow L_l \Phi) + \Gamma(N_i \rightarrow L_l^C \Phi^l) \right] / H(T = m_{N_i}),
\]

where \( \Gamma \) is the decay rate of a heavy Majorana neutrino mass eigenstate \( N_i \) with mass \( m_{N_i} \).

Successful leptogenesis with hierarchical heavy neutrinos needs a relatively small value of \( K_{N_i} \), typically \( K_{N_i} \sim 1 \). A large CP-asymmetry allows successful resonant leptogenesis with very large values of \( K_{N_i} \sim 1000 \) (a very small departure from thermal equilibrium). Conditions close to thermal equilibrium remove the dependence of the baryon asymmetry on the initial heavy neutrino abundance or any initial lepton or baryon asymmetries.

In general, the matrix of neutrino Yukawa couplings \( h^\nu \) will have different entries coupling a given heavy neutrino to each lepton flavour. This should be taken into account by considering the dynamical generation of each lepton flavour asymmetry. In a minimal model, an asymmetry in just one flavour is needed to produce the baryon asymmetry of the Universe and this asymmetry could be produced via the out of equilibrium decay of just one heavy Majorana neutrino.\(^6\) This is possible because \( B + L \) violating sphaleron interactions individually preserve \( \frac{4}{3} B - L_{e,\mu,\tau} \), and will therefore convert an individual flavour asymmetry into a baryon asymmetry.\(^5\)

These requirements mean that \( K_{N_i} \) should be relatively small for only 1 heavy neutrino \( N_i \), and the \( K_{N_i}^l \) for only one lepton flavour \( l \) needs to be small to avoid washing out the produced asymmetry. Thus the requirement of small Yukawa couplings, coming from demanding successful thermal leptogenesis, can be substantially relaxed.

Once the model is specified, to accurately determine the baryon asymmetry one must solve the relevant Boltzmann equations. It is possible to determine and solve coupled Boltzmann equations for each lepton flavour asymmetry and also the baryon asymmetry.\(^3\) Solving a separate Boltzmann equation for the baryon asymmetry allows the inclusion of effects due to the sphaleron mediated processes dropping out of equilibrium at temperatures less than the critical temperature for electroweak symmetry breaking.

If \( K_{N_i} \gtrsim 1 \), an order of magnitude estimate of the baryon asymmetry may be obtained using

\[
\eta_B \sim 10^{-2} \sum_{l=e,\mu,\tau} \sum_{i=1}^3 \delta_{N_i}^l \delta_{N_i}^l K_{N_i}^l / K_{N_i}^l,
\]

where \( \delta_{N_i}^l \) is the CP-asymmetry in the decay of \( N_i \) into leptons of flavour \( l \), \( K_{N_i} = \sum_l K_{N_i}^l \) and \( K_l = \sum_i e^{-(m_{N_i} - m_{N_i})/m_{N_i}} K_{N_i}^l \). It can be shown that the inclusion of the dynamical effects of lepton flavours can alter the predicted baryon asymmetry in some resonant leptogenesis scenarios by factors of order \( 10^6 \). Even in scenarios with a mildly hierarchical spectrum of heavy neutrinos the effects of individual lepton flavours can alter predictions by a factor of 10.

3 A Model for Resonant \( \tau \)-Leptogenesis

To fully exploit the phenomenological possibilities offered by low-scale, resonant leptogenesis it is necessary to abandon the usual see-saw paradigm that small light neutrino masses arise because of the large Majorana masses of gauge singlet right-handed neutrinos. Instead, we will consider models where small light neutrino masses arise due to approximate flavour symmetries. Approximate flavour symmetries can also motivate a nearly degenerate spectrum of heavy Majorana neutrino masses.

As a definite example,\(^6\) we will consider a model with a heavy neutrino sector that is SO(3) symmetric in the absence of Yukawa couplings.\(^8\) This will provide 3 nearly degenerate heavy

\(^6\)Several other works have studied flavour issues and leptogenesis.\(^7\) In the low scale \( T < 10^6 \) GeV scenarios discussed here, all the charged lepton Yukawa couplings are in equilibrium, favouring the basis in which they are diagonal. The effects of quantum oscillations amongst the flavour asymmetries are therefore suppressed.
Majorana neutrinos. Specifically, the $\nu_{iR}$ transform as triplets under this symmetry, whilst all other fields are singlets. Therefore, $M_S$ can be written $M_S = m_N 1_3 + \Delta M_S$, where $\Delta M_S$ parameterizes deviations from the imposed $SO(3)$ symmetry.

The presence of the charged lepton Yukawa couplings breaks the leptonic flavour symmetry of the model down to the usual three $U(1)_l$ flavour symmetries of the SM and the $SO(3)$ symmetry of the gauge singlet neutrinos. Finally, when imposing the neutrino Yukawa couplings we choose to leave a particular $SO(2)\cong U(1)_l$ subgroup of the total flavour symmetry unbroken (to leading order). Specifically, all lepton doublets are coupled to the particular heavy neutrino combination: $\frac{1}{\sqrt{2}}(\nu_{2R} + i\nu_{3R})$. The $U(1)_l$ charges of the fields are therefore: $Q(L_1) = Q(l_{iL}) = 1$, $Q(\frac{1}{\sqrt{2}}(\nu_{2R} + i\nu_{3R})) = -Q(\frac{1}{\sqrt{2}}(\nu_{2R} - i\nu_{3R})) = 1$, $Q(\nu_{1R}) = 0$. This symmetry leads to the following generic structure for the neutrino Yukawa couplings:

$$H^{\nu R} = \begin{pmatrix} 0 & a e^{-i\pi/4} & a e^{i\pi/4} \\ 0 & b e^{-i\pi/4} & b e^{i\pi/4} \\ 0 & c e^{-i\pi/4} & c e^{i\pi/4} \end{pmatrix} + \begin{pmatrix} \varepsilon_e & 0 & 0 \\ \varepsilon_\mu & 0 & 0 \\ \varepsilon_\tau & 0 & 0 \end{pmatrix}. \tag{3}$$

The $U(1)_l$ symmetry is broken by both $\Delta M_S$ and the second term in equation (3)\footnote{We will not consider the origin of the symmetry breaking terms here, but they could arise for example in a Froggatt-Nielsen scenario. Fine tuning is a concern, but in the examples considered here we have imposed a `naturalness criterion' requiring that the cancellation between tree and radiative effects is less than 1 part in $\sim 20$.}

The parameters $a$, $b$ and $c$ in equation (3) are arbitrary complex parameters which should be smaller than $10^{-2}$ to provide a light neutrino sector compatible with current data. In addition, they allow the choice of a specific flavour direction in which the lepton asymmetry can be preserved. For example, the requirement that an excess in $L_\tau$ be protected from wash-out effects leads to the constraint $|c| \lesssim 10^{-5}$. The small $U(1)_l$ breaking parameters $\varepsilon_{e,\mu,\tau}$ should be of order $10^{-6}$ allowing small light neutrino masses and one heavy neutrino to decay sufficiently far out of thermal equilibrium. Note that the neutrino Yukawa couplings in this model span a range comparable to that of the charged lepton Yukawa couplings.

Figure 1 shows the results of solving the Boltzmann equations with various initial conditions in a scenario with $m_N = 250$ GeV, in complete accordance with light neutrino data and an inverted hierarchical light neutrino spectrum\footnote{The Boltzmann equations were solved using the Fortran code LeptoGen, available from http://www.ippp.dur.ac.uk/~teju/leptogen.}. The final baryon asymmetry is almost completely independent of the initial baryon and lepton asymmetries and the initial abundance of heavy neutrinos. In this scenario, $|a| \sim |b| \sim 10^{-2}, |c| \sim 10^{-6}, |\varepsilon_{e,\mu,\tau}| \sim 10^{-6}$ and $(\Delta M_S)_{ij}/m_N$ ranged between $10^{-5}$ and $10^{-9}$.
This choice of parameters leads to large $K_{N_{12}}^{e,\mu} \sim 10^{10}$, whilst $K_{N_3}^{\tau} \sim 10$ is relatively small, allowing $N_3$ to decay sufficiently out of thermal equilibrium to produce a lepton asymmetry. A $\tau$-lepton asymmetry is protected from washout by $K_{N_3}^{\tau} \sim 5$. With these $K$-factors a CP-asymmetry $\delta_{N_3} \sim -10^{-6}$ is sufficient to generate the observed baryon asymmetry.

4 Phenomenology

The observation of lepton number violation in neutrinoless double beta decay (0$\nu$\beta\beta-decay) would be an important piece of evidence pointing to leptogenesis as the origin of the baryon asymmetry of the Universe. In the example scenario studied, the effective Majorana mass is $|\langle m \rangle| \sim 0.013$ eV. This figure is close to the sensitivity of proposals for future 0$\nu$\beta\beta-decay experiments, which should at least significantly constrain the parameter space for these models.

Heavy Majorana neutrinos with relatively large Yukawa couplings may induce lepton flavour violating (LFV) couplings to the photon and Z-boson. In a scenario with large Yukawa couplings to the electron and muon the LFV decays $\mu \to e\gamma$ and $\mu \to eee$ can be induced. The same LFV couplings can give rise to the coherent conversion of $\mu \to e$ in nuclei e.g. $\mu^- \to^{48}_{22} Ti \to e^-^{48}_{22} Ti$.

In the previous example, the branching fraction for the decay $\mu \to e\gamma$ is given by

$$B(\mu \to e\gamma) = 9 \times 10^{-4} |a|^2 |b|^2 v^4/m_N^4,$$

where $v \simeq 246$ GeV is the vacuum expectation value of the Higgs field. For $m_N = 250$ GeV we find $B(\mu \to e\gamma) \simeq 10^{-12}$, which is well within reach of the experiment proposed by the MEG collaboration.\(^9\) The branching fractions for $\mu \to eee$ and $\mu \to e$ conversion in $^{48}_{22} Ti$ are related to this result, and for $m_N = 250$ GeV one obtains

$$B(\mu \to eee) \simeq 1.4 \times 10^{-2} \times B(\mu \to e\gamma) \simeq 1.4 \times 10^{-14},$$
$$B_{^{48}_{22} Ti}(\mu \to e) \simeq 0.46 \times B(\mu \to e\gamma) \simeq 4.5 \times 10^{-13}.$$  

The result for $B_{^{48}_{22} Ti}(\mu \to e)$ falls well within reach of the experiments proposed by the MECO and PRIME collaborations.\(^10\)

Finally, in the previous example, large couplings between $N_{1,2}$ and $e, \mu$ mean that it should be possible to produce $N_{1,2}$ at a future $e^+e^-$ or $\mu^+\mu^-$ linear collider. Taking backgrounds into account, the study\(^11\) finds that an $e^+e^-$ linear collider $\sqrt{s} = 0.5$ TeV with would be sensitive to $|a|v/m_N \sim 0.7 \times 10^{-2}$.

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