

# The Little Review on Leptogenesis

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**Abstract.** This is a brief review on the scenario of baryogenesis through leptogenesis. Leptogenesis is an appealing scenario that may relate the observed baryon asymmetry in the Universe to the low-energy neutrino data. In this review talk, particular emphasis is put on recent developments on the field, such as the flavourdynamics of leptogenesis and resonant leptogenesis near the electroweak phase transition. It is illustrated how these recent developments enable the modelling of phenomenologically predictive scenarios that can directly be tested at the LHC and indirectly in low-energy experiments of lepton-number and lepton-flavour violation.

## 1. Introduction

An elegant framework to consistently address the observed Baryon Asymmetry in the Universe (BAU) [1, 2] is leptogenesis [3]. According to the standard paradigm of leptogenesis [4], minimal extensions of the Standard Model (SM) realize heavy Majorana neutrinos of masses related to the Grand Unified Theory (GUT) scale  $M_{\text{GUT}} \sim 10^{16}$  that decay out of equilibrium and create a net excess of lepton number ( $L$ ). This excess in  $L$  gets then reprocessed into the observed baryon number ( $B$ ), through the  $(B + L)$ -violating sphaleron interactions [5]. The attractive feature of such a scenario is that the GUT-scale heavy Majorana neutrinos could also explain the observed smallness in mass of the SM light neutrinos by means of the so-called seesaw mechanism [6].

The original GUT-scale leptogenesis scenario, however, runs into certain difficulties, when one attempts to explain the flatness of the Universe and other cosmological data [2] within supergravity models of inflation. To avoid overproduction of gravitinos  $\tilde{G}$  whose late decays may ruin the successful predictions of Big Bang Nucleosynthesis (BBN), the reheat temperature  $T_{\text{reh}}$  of the Universe should be lower than  $10^9$ – $10^6$  GeV, for  $m_{\tilde{G}} = 8$ – $0.2$  TeV [7]. This implies that the heavy Majorana neutrinos should accordingly have masses as low as  $T_{\text{reh}} \lesssim 10^9$  GeV, thereby rendering the relation of these particles with GUT-scale physics less natural. On the other hand, it proves very difficult to directly probe the heavy-neutrino sector of such a model at high-energy colliders, e.g. at the LHC or ILC, or in any other foreseeable experiment.

A potentially interesting solution to the above problems may be obtained within the framework of resonant leptogenesis (RL) [8]. The key aspect of RL is that self-energy effects dominate the leptonic asymmetries [9], when two heavy Majorana neutrinos happen to have

a small mass difference with respect to their actual masses. If this mass difference becomes comparable to the heavy neutrino widths, a resonant enhancement of the leptonic asymmetries takes place that may reach values  $\mathcal{O}(1)$  [8, 10]. An indispensable feature of RL models is that flavour effects due to the light-to-heavy neutrino Yukawa couplings [11, 12, 13] play a dramatic role and can modify the predictions for the BAU by many orders of magnitude [11, 12]. Most importantly, these flavour effects enable the modelling [12] of minimal RL scenarios with electroweak-scale heavy Majorana neutrinos that could be tested at the LHC [14, 15, 16] and in other non-accelerator experiments, while maintaining agreement with the low-energy neutrino data. Many variants of RL have been proposed in the literature [17, 18], including soft leptogenesis [19] and radiative leptogenesis [20].

In this review we give a brief exposition of selected topics related to recent advancements of the field. In detail, in Section 2, we remind ourselves about Sakharov’s basic conditions for baryogenesis and list a few classical scenarios that have been suggested in the literature, including leptogenesis. In Section 3 we discuss the importance of the various sources of flavour effects on leptogenesis. Section 4 presents the field theory of RL, including a discussion of its flavourdynamics. Section 5 discusses the potential particle-physics implications of RL. Finally, Section 6 summarizes the main points of this review.

## 2. Matter–Antimatter Asymmetry and Leptogenesis

There are two pieces of observational evidence that establish the domination of the observable matter over antimatter in our Universe. The first stems from the analysis of the power spectrum of the cosmic microwave background [2]. In particular, the baryon-to-photon ratio  $\eta_B$  of number densities is found to be

$$\eta_B^{\text{CMB}} = \frac{n_B}{n_\gamma} = 6.1_{-0.2}^{+0.3} \times 10^{-10} . \quad (1)$$

The second piece of evidence originates from the concordance of the light elements and big bang nucleosynthesis, which gives [21]

$$\eta_B^{\text{BBN}} = 3.4\text{--}6.9 \times 10^{-10} . \quad (2)$$

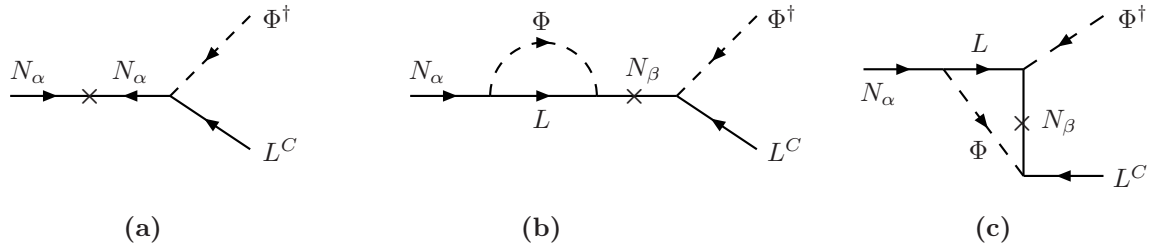
In spite of their different origin, the two deduced values for  $\eta_B$  are in good agreement with one another.

More than fifty years ago [1], Sakharov has studied the conditions for which a baryon asymmetry may successfully be generated. These basic conditions are:

- the theory should have B-violating interactions.
- the interactions should violate both C and CP.
- the processes of a net baryon-number generation should have a degree of irreversibility due to an out-of-thermal equilibrium dynamics.

There are many scenarios that have been proposed in the existing literature. A few classical examples are:

- *Baryogenesis through the decay of a heavy particle.* This scenario is based on the out-of-equilibrium B-violating decay of a heavy GUT particle, e.g. in SO(10) theory [22]. A serious constraint difficult to overcome in these scenarios comes from the strict experimental lower limit on proton’s long lifetime [21].
- *Baryogenesis at the electroweak phase transition.* In this scenario, BAU is generated by (B + L)-violating sphaleron interactions at  $T \sim T_c \approx 140$  GeV, through a first order phase



**Figure 1.** Feynman diagrams contributing to the  $L$ -violating decays of heavy Majorana neutrinos,  $N_\alpha \rightarrow L^C \Phi^\dagger$ , where  $L$  and  $\Phi$  represent lepton and Higgs-boson iso-doublets, respectively: (a) tree-level graph, and (b) self-energy and (c) vertex graph.

transition [5]. However, such a scenario cannot be realized within the Standard Model (SM) [23], since the phase transition is not sufficiently strong enough given the direct limits on the Higgs-boson mass. In this case, one needs to resort to extended models with extended Higgs sectors, such as the MSSM [24].

- *Baryogenesis through leptogenesis.* This scenario combines the two key ideas of the above two scenarios. Leptogenesis is based on the out-of-equilibrium  $L$ -violating decays of heavy Majorana neutrinos. These decays produce a net lepton asymmetry, which is converted into the observed BAU through  $(B + L)$ -violating sphaleron interactions [3].

Several variants of leptogenesis have been discussed so far in the literature [25]. These include Dirac leptogenesis [26], non-thermal leptogenesis [27], Affleck-Dine and spontaneous leptogenesis [28], and RL that will be discussed more extensively in Section 4.

### 3. The Flavourdynamics of Leptogenesis

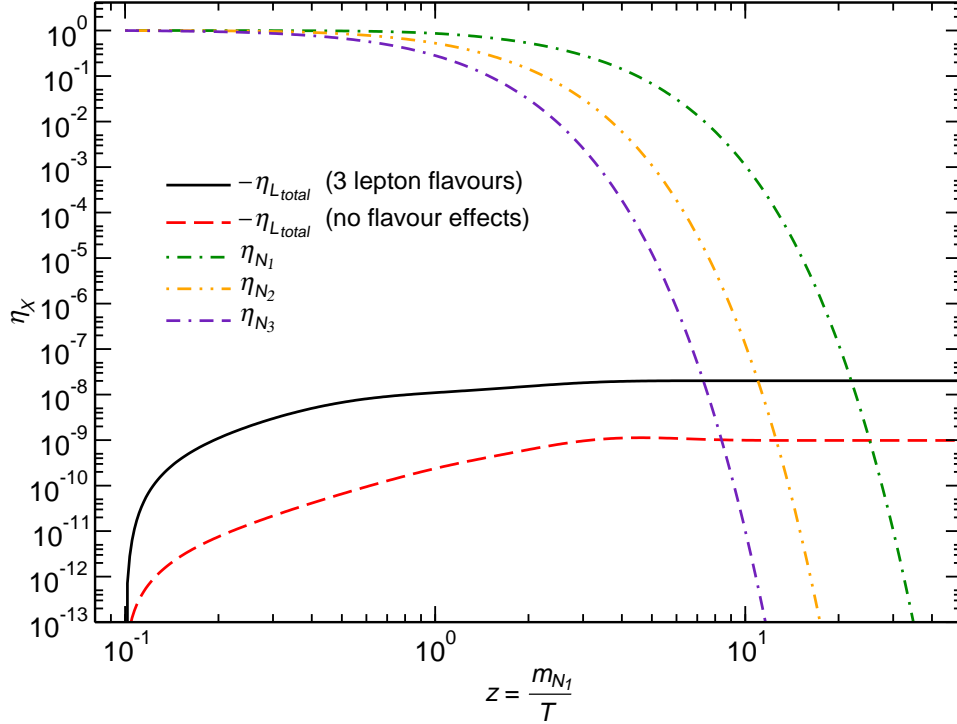
An important recent development is the understanding of the special role that flavour effects can play in leptogenesis. In fact, this special role of flavour effects on scenarios of baryogenesis is not new [29] and has already been recognised in the past [30, 31]. The key observation here is that sphalerons do not only conserve  $B - L$ , but also the individual quantum numbers:

$$\frac{1}{3}B - L_{e,\mu,\tau}.$$

Thus, a generated baryon asymmetry can be protected from erasure if a particular lepton flavour combination or an individual lepton number  $L_{e,\mu,\tau}$  is out of thermal equilibrium. For instance, this idea has been used in [30] to show how specific R-parity-violating supersymmetric scenarios do not wash out any baryon asymmetry that could eventually account for the observed matter-antimatter asymmetry in the Universe.

The new aspect of flavourdynamics in leptogenesis is that the BAU can now be both *generated from* and *protected in a single* lepton flavour [11], e.g. in the  $\tau$ -lepton number  $L_\tau$ . In order to better understand this, we need first to clarify the two sources of flavour effects:

- *flavour effects due to charged-lepton Yukawa couplings  $h_{e,\mu,\tau}$ .* These effects are related to the interactions mediated by charged-lepton Yukawa couplings [32]. If these interactions are in or out of thermal equilibrium at the scale, at which leptogenesis takes place, then the predicted value for the BAU gets modified significantly, typically up to one order of

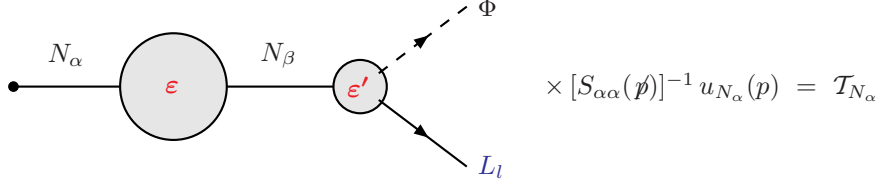


**Figure 2.** Numerical estimates for a scenario with  $m_{N_1} = 10^{10}$  GeV,  $m_{N_2} = 2m_{N_1}$ ,  $m_{N_3} = 3m_{N_1}$ .

magnitude [33]. Thus, if the leptogenesis scale is higher than  $10^9$  TeV, then all charged-lepton Yukawa interactions are out of thermal equilibrium, whereas if the scale is below 10 TeV, all charged-lepton Yukawa couplings are in thermal equilibrium [31] and their full impact on the BAU needs be taken into account [12].

- *flavour effects due to heavy-neutrino Yukawa couplings  $h_{l\alpha}^\nu$ .* These effects are governed by heavy-neutrino Yukawa-coupling interactions [11, 13]. In this case, one may consider scenarios that would generate an excess in an individual lepton number, e.g.  $L_\tau$ , and protect the resulting asymmetry from potentially large wash-out effects due to sphalerons. For example, one could arrange the neutrino Yukawa couplings  $h_{l\alpha}^\nu$ , such that all heavy Majorana neutrinos  $N_{1,2,3}$  decay out of equilibrium in  $L_\tau$ ,  $N_1$  decays out of equilibrium in  $L_{e,\mu,\tau}$ , but  $N_{2,3}$  could decay very rapidly into  $L_{e,\mu}$ . In Fig. 2, we give predictions for an hierarchical scenario with and without neutrino Yukawa coupling effects included [12, 13]. The two predictions can differ by up to one order of magnitude. As we will discuss in Section 4, however, flavour effects due to  $h_{l\alpha}^\nu$  may modify the BAU predictions by many orders of magnitude, e.g. bigger than  $10^6$ , in RL models, thereby rendering their inclusion indispensable [11, 12].

We close this section by mentioning that flavour effects do also play an important role in collision terms of the Boltzmann equations that describe scatterings [10, 12]. These are  $\Delta L = 1$  Yukawa and gauge-mediated scatterings, as well as  $\Delta L = 0, 2$  scatterings mediated by heavy Majorana neutrinos.



**Figure 3.** LSZ-type approach to mixing and decay of heavy Majorana neutrinos

#### 4. Resonant Leptogenesis

A scenario inspired by the dynamics of the  $K^0\bar{K}^0$  system [34] is resonant leptogenesis [8, 10]. Resonant leptogenesis is based on the observation that self-energies dominate the lepton asymmetries if  $|m_{N_1} - m_{N_2}| \ll m_{N_{1,2}}$  [9]. However, as we will see below, the proper inclusion of heavy-neutrino width effects is necessary in order to obtain a non-singular well-behaved analytic expression for the leptonic asymmetries.

There are several variants of RL:

- *Soft RL* [19]. In this scenario, leptogenesis results from sneutrino decays in the MSSM. CP-violating soft SUSY-breaking parameters may split the degeneracy within a single generation of right-handed sneutrino states, which leads to a resonant enhancement of the leptonic asymmetries.
- *Radiative RL* [20]. In this case, the heavy Majorana neutrinos are exactly degenerate at the GUT scale. Then, small mass differences among the heavy Majorana neutrino states are generated via renormalization group (RG) running from the GUT scale down to the leptogenesis scale. These small mass differences give rise to RL.
- *Coherent RL (via sterile neutrino oscillations)* [35, 36]. This scenario relies on the possibility that sterile neutrinos with masses well below the critical temperature of the electroweak phase transition, e.g. of order 10 GeV, maintain the coherence of their CP asymmetric oscillations and can thus produce an enhanced leptonic asymmetry of the amount needed to create the observed BAU.

We will now briefly review the field theory developed in [8, 10] for thermal RL. This is based on an effective LSZ-type formalism [37] with the aim to incorporate both the mixing and the decay of the heavy Majorana neutrinos. As shown in 3, one needs to resum self-energy graphs, whose absorptive parts regularize the tree-level mass singularity when  $m_{N_\alpha} \rightarrow m_{N_\beta}$ . For a two-generation heavy-neutrino model, the resummed propagator matrix reads:

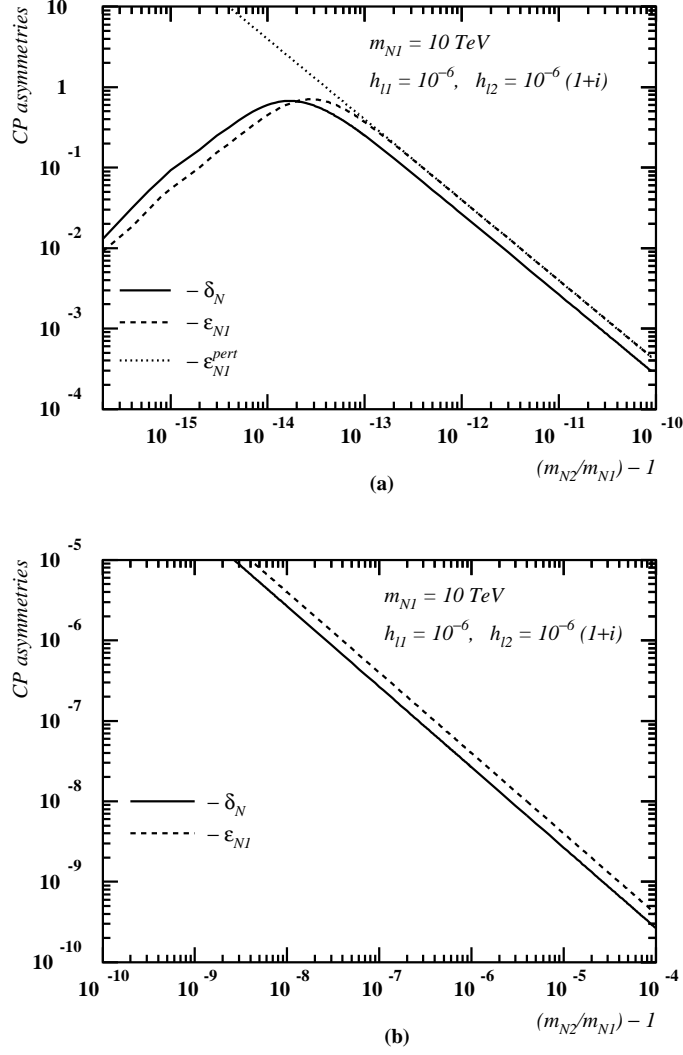
$$S_{\alpha\beta}(p) = \begin{pmatrix} p - m_{N_1} + \Sigma_{11}(p) & \Sigma_{12}(p) \\ \Sigma_{21}(p) & p - m_{N_2} + \Sigma_{22}(p) \end{pmatrix}^{-1}. \quad (3)$$

Extensive technical details of this formalism may be found in [8, 10].

Making use of the corresponding  $K^0\bar{K}^0$  terminology, there are two types of CP violation in the leptonic decays of the heavy Majorana neutrinos:

- $\varepsilon'$ -type CP violation. In this case, the required CP violation arises from the interference of the tree-level graph with the absorptive of the self-energy effects.

$$\varepsilon'_{N_\alpha} = \frac{\text{Im}(h^{\nu\dagger} h^\nu)_{\alpha\beta}^2}{(h^{\nu\dagger} h^\nu)_{\alpha\alpha} (h^{\nu\dagger} h^\nu)_{\beta\beta}} \left( \frac{\Gamma_{N_\beta}}{m_{N_\beta}} \right) f \left( \frac{m_{N_\beta}^2}{m_{N_\alpha}^2} \right), \quad (4)$$



**Figure 4.**  $\varepsilon$ - and  $\varepsilon'$ -type contributions to the leptonic asymmetries in a model with two heavy Majorana neutrinos, with  $\delta_{N_{1,2}} \approx \varepsilon_{N_{1,2}} + \varepsilon'_{N_{1,2}}$  and  $\delta_N \approx \delta_{N_1} + \delta_{N_2}$ .

where

$$\Gamma_{N_\beta} = \frac{(h^{\nu\dagger} h^\nu)_{\beta\beta}}{8\pi} m_{N_\beta}$$

is the tree-level decay width of  $N_\beta$ , and  $f(x) = \sqrt{x}[1 - (1+x)\ln(1+1/x)]$  is the Fukugita–Yanagida loop function [3].

- $\varepsilon$ -type CP violation.

$$\varepsilon_{N_\alpha} = \frac{\text{Im}(h^{\nu\dagger} h^\nu)_{\alpha\beta}^2}{(h^{\nu\dagger} h^\nu)_{\alpha\alpha} (h^{\nu\dagger} h^\nu)_{\beta\beta}} \left( \frac{\Gamma_{N_\beta}}{m_{N_\beta}} \right) \frac{(m_{N_\alpha}^2 - m_{N_\beta}^2) m_{N_\alpha} m_{N_\beta}}{(m_{N_\alpha}^2 - m_{N_\beta}^2)^2 + m_{N_\alpha}^2 \Gamma_{N_\beta}^2} \quad (5)$$

Note that  $\varepsilon_{N_\alpha}$  dominate over the  $\varepsilon'_{N_\alpha}$  when  $m_{N_\alpha} - m_{N_\beta} \ll m_{N_{\alpha,\beta}}$ . In the limit  $m_{N_\alpha} \rightarrow m_{N_\beta}$ , the would-be singularity is regularized by the decay width of the heavy neutrino  $N_\beta$ .

From (5), the following *two* conditions for  $\mathcal{O}(1)$  leptonic asymmetries may be obtained [8]:

$$(i) \quad m_{N_\alpha} - m_{N_\beta} \sim \frac{1}{2} \Gamma_{N_{\alpha,\beta}}, \quad (ii) \quad \frac{\text{Im}(h^{\nu\dagger} h^\nu)_{\alpha\beta}^2}{(h^{\nu\dagger} h^\nu)_{\alpha\alpha} (h^{\nu\dagger} h^\nu)_{\beta\beta}} \sim 1. \quad (6)$$

Figure 4 shows the impact of the above two conditions on the leptonic asymmetries for a low-scale seesaw model. The generic feature of the resonant conditions remains valid even in the presence of flavour effects [12].

As was already mentioned in Section 3, flavour effects due to neutrino Yukawa couplings play an important role in models of RL. As an illuminating example, we will consider a scenario of RL, which uses the  $\tau$ -lepton number to generate the BAU and has the potential to be directly tested as we will see in the next section. In such a model, the smallness of the light neutrinos is not due to the usual seesaw mechanism [6], but thanks to the approximate breaking of lepton flavour symmetries [38, 39, 14, 11]. In detail, the heavy neutrino sector is initially assumed to be  $\text{SO}(3)$ , and gets broken down to the identity  $\mathbf{I}$  via the neutrino Yukawa couplings in two steps [11]:

$$\text{SO}(3) \xrightarrow{\sim h_\tau} \text{SO}(2) \simeq \text{U}(1)_l \xrightarrow{\sim h_\xi} \mathbf{I}.$$

The neutrino Dirac mass matrix that results from the Yukawa sector reads:

$$m_D = \frac{v}{\sqrt{2}} \begin{pmatrix} \varepsilon_e & a e^{-i\pi/4} & a e^{i\pi/4} \\ \varepsilon_\mu & b e^{-i\pi/4} & b e^{i\pi/4} \\ \varepsilon_\tau & c e^{-i\pi/4} & c e^{i\pi/4} \end{pmatrix}, \quad (7)$$

with  $a \sim b \sim 10^{-2} \sim h_\tau$  and  $c, |\varepsilon_l| \sim 10^{-7} \sim h_e$ . Likewise, the heavy Majorana mass matrix  $m_M$  deviates from  $\mathbf{I}$  by terms  $\Delta m_M \sim (h^{\nu T} h^{\nu*} + h^{\nu\dagger} h^\nu) m_N$ , as naively expected by RG effects, where  $m_N$  is an  $\text{SO}(3)$  symmetric Majorana mass in the range 0.1–1 TeV. Specifically, the light neutrino mass matrix is given by [12]

$$m_\nu^{\text{light}} = \frac{v^2}{2m_N} \begin{pmatrix} \frac{\Delta m_N}{m_N} a^2 - \varepsilon_e^2 & \frac{\Delta m_N}{m_N} ab - \varepsilon_e \varepsilon_\mu & \frac{\Delta m_N}{m_N} ac - \varepsilon_e \varepsilon_\tau \\ \frac{\Delta m_N}{m_N} ab - \varepsilon_e \varepsilon_\mu & \frac{\Delta m_N}{m_N} b^2 - \varepsilon_\mu^2 & \frac{\Delta m_N}{m_N} bc - \varepsilon_\mu \varepsilon_\tau \\ \frac{\Delta m_N}{m_N} ac - \varepsilon_e \varepsilon_\tau & \frac{\Delta m_N}{m_N} bc - \varepsilon_\mu \varepsilon_\tau & \frac{\Delta m_N}{m_N} c^2 - \varepsilon_\tau^2 \end{pmatrix},$$

where

$$\Delta m_N = 2(\Delta m_M)_{23} + i[(\Delta m_M)_{33} - (\Delta m_M)_{22}], \quad \frac{b}{a} = \frac{19}{50},$$

and (in  $\sim 10^{-7}$  units)

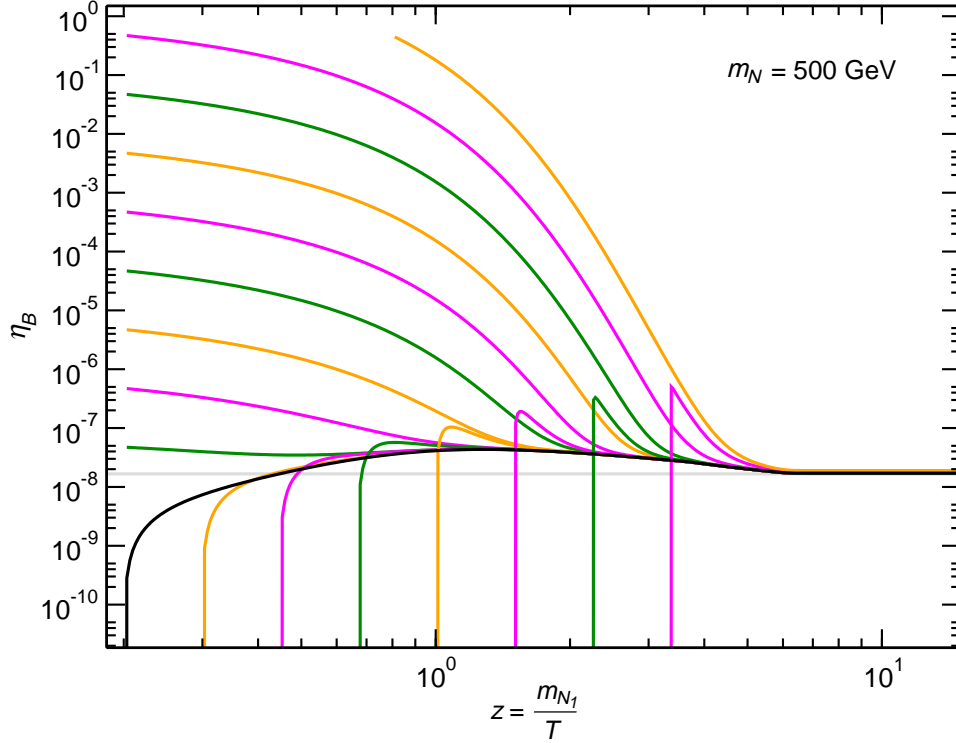
$$\sqrt{\frac{\Delta m_N}{m_N}} a = 2, \quad \varepsilon_e = 2 + \frac{21}{250}, \quad \varepsilon_\mu = \frac{13}{50}, \quad \varepsilon_\tau = -\frac{49}{128}.$$

The above choice of parameters predicts a light-neutrino sector that realizes an inverted mass hierarchy,  $m_{\nu_3} < m_{\nu_1} < m_{\nu_2}$ , with

$$\begin{aligned} m_{\nu_2}^2 - m_{\nu_1}^2 &= 7.54 \times 10^{-5} \text{ eV}^2, & m_{\nu_1}^2 - m_{\nu_3}^2 &= 2.45 \times 10^{-3} \text{ eV}^2, \\ \sin^2 \theta_{12} &= 0.362, & \sin^2 \theta_{23} &= 0.341, & \sin^2 \theta_{13} &= 0.047. \end{aligned}$$

The above is compatible with low-energy neutrino oscillation data at the  $3\sigma$  level [40].

In Fig. 5, we give numerical estimates of the baryon asymmetry  $\eta_B$  for different initial conditions in the resonant  $\tau$ -leptogenesis ( $\text{R}\tau\text{L}$ ) model described above. The heavy Majorana



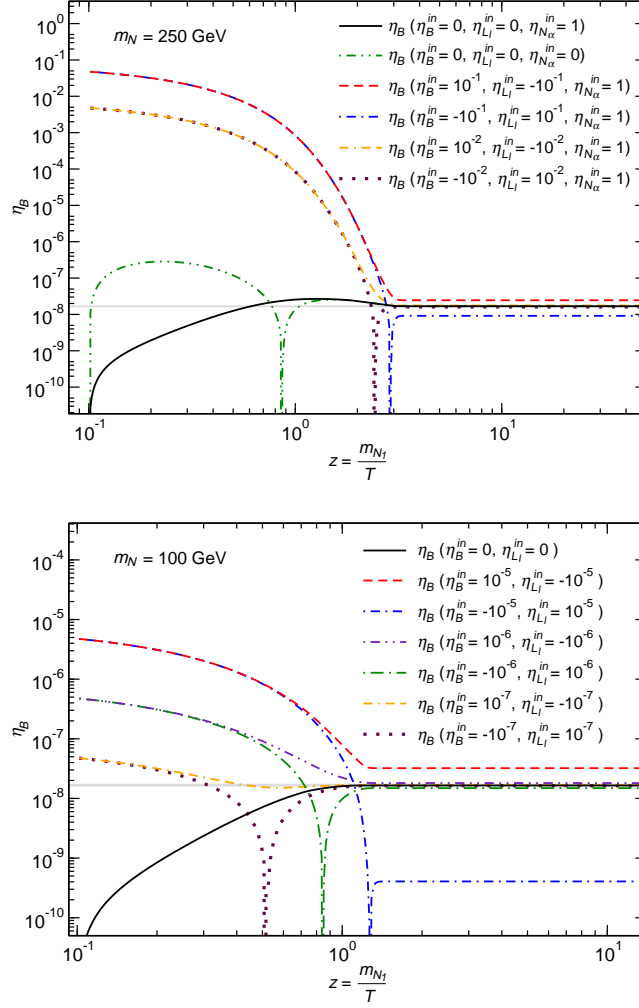
**Figure 5.** Numerical estimates of the baryon asymmetry  $\eta_B$  for different initial conditions in the  $R\tau L$  model with  $m_N = 500$  GeV. The horizontal line indicates the observed baryon asymmetry as required just after the sphalerons freeze out.

mass scale is  $m_N = 500$  GeV. Our results show that the baryon asymmetry  $\eta_B$  is almost independent on the initial conditions. This property is a consequence of the fixed-point dynamics which exhibits a system that is in quasi-thermal equilibrium. As can be seen from the upper panel in Fig. 6, this almost independence of  $\eta_B$  on the initial conditions persists, even if the heavy Majorana mass  $m_N$  is as low as 250 GeV. For lower masses, e.g.  $m_N = 100$  GeV (see lower panel in Fig. 6), this feature becomes less pronounced, and the dynamics of leptogenesis at the electroweak phase transition becomes more involved. In this case, one needs to take into consideration contributions to the leptonic asymmetries that originate from the transverse polarizations of the  $W$ - and  $Z$ -boson gauge fields [41].

It is therefore tempting to introduce some measure that could quantify the initial condition dependence of the different models suggested in the literature. To this end, we may define a  $Q$ -factor

$$Q = \ln \left| \frac{\delta\eta_B^{\text{in}}}{\delta\eta_B^{\text{fin}}} \right|, \quad (8)$$

where  $T_{\text{in}}$  is a typical initial temperature when the baryogenesis mechanism becomes active. For instance, for leptogenesis, this could be identified as  $T_{\text{in}} \sim m_N$  or for electroweak baryogenesis, one may assume that  $T_{\text{in}}$  is the critical temperature of the electroweak phase transition  $T_c^{\text{EW}}$ . Correspondingly,  $T_{\text{fin}}$  is the freeze-out temperature of the baryogenesis mechanism under study. In terms of the  $Q$ -factor, one could quantify the actual dependence of the model for baryogenesis on the initial conditions. Thus, if  $Q \gg 1$ , this would imply that there is very weak dependence of the BAU on the initial conditions, whereas for  $Q < 0$  this would signify that the BAU does



**Figure 6.**  $\eta_B$  versus  $z = m_{N_1}/T$  in the R $\tau$ L model with  $m_N = 250$  GeV (upper panel) and  $m_N = 100$  GeV (lower panel).

strongly depend on these. Evidently, RL models, as the one analyzed here, have large values of  $Q$  and so render them quite predictive.

## 5. Particle-Physics Phenomenology of Resonant Leptogenesis

RL models can give rise to a number of phenomenologically testable signatures. Here, we will present the generic predictions of R $\tau$ L models for the  $0\nu\beta\beta$  decay, and for processes of lepton flavour violation (LFV), such as  $\mu \rightarrow e\gamma$ ,  $\mu \rightarrow eee$  and  $\mu \rightarrow e$  conversion in nuclei. Finally, we give numerical estimates of production cross sections of heavy Majorana neutrinos at the LHC.

### 5.1. $0\nu\beta\beta$ Decay

Neutrinoless double beta decay ( $0\nu\beta\beta$ ) corresponds to a process in which two single  $\beta$  decays [42] occur simultaneously in one nucleus. As a consequence of this, a nucleus ( $Z, A$ ) gets converted into a nucleus ( $Z + 2, A$ ), i.e.

$${}^A_Z X \rightarrow {}^A_{Z+2} X + 2e^-.$$

Evidently, this process violates  $L$ -number by two units and can naturally take place in minimal RL models, in which the observed light neutrinos are Majorana particles. The observation of such a process would provide further information on the structure of the light neutrino mass matrix  $\mathbf{m}^\nu$ . In particular, the transition rate of  $0\nu\beta\beta$  decay is governed by the effective Majorana neutrino mass  $\langle m \rangle$ . For the R $\tau$ L model, the effective neutrino mass is given by

$$|\langle m \rangle| = |(\mathbf{m}^\nu)_{ee}| = \frac{v^2}{2m_N} \left| \frac{\Delta m_N}{m_N} a^2 - \varepsilon_e^2 \right| \approx 0.013 \text{ eV} . . \quad (9)$$

Such a prediction lies within the reach of future experiments which will be sensitive to values of  $|\langle m \rangle|$  of order  $10^{-2}$ , such as SuperNEMO [43].

### 5.2. $\mu \rightarrow e\gamma$ , $\mu \rightarrow eee$ and $\mu \rightarrow e$ conversion

Quantum effects due to heavy Majorana neutrinos may induce LFV couplings to the photon ( $\gamma$ ) and the  $Z$  boson. These couplings give rise to LFV decays, such as  $\mu \rightarrow e\gamma$  [44],  $\mu \rightarrow eee$  [45] and  $\mu \rightarrow e$  conversion [46]. More details on the actual calculation of these processes may be found in [45, 47]. Also, related phenomenological analyses of LFV effects in the SM with singlet neutrinos may be found in [48, 49, 50, 51].

Let us first consider the decay  $\mu \rightarrow e\gamma$ . In the R $\tau$ L model, the branching fraction for  $\mu \rightarrow e\gamma$  is given by

$$B(\mu \rightarrow e\gamma) = 9 \cdot 10^{-4} \times \frac{|a|^2 |b|^2 v^4}{m_N^4} . \quad (10)$$

The theoretical prediction should be contrasted with the experimental upper limit:

$$B_{\text{exp}}(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11} . \quad (11)$$

For the particular scenario considered in Section 4, we find  $B(\mu \rightarrow e\gamma) \sim 10^{-12}$ . These values are well within reach of the MEG collaboration, which will be sensitive to  $B(\mu \rightarrow e\gamma) \sim 10^{-14}$  [52].

We then consider the decay  $\mu \rightarrow eee$ . Its branching ratio  $B(\mu \rightarrow eee)$  may be related to  $B(\mu \rightarrow e\gamma)$  through [12]:

$$B(\mu \rightarrow eee) \simeq 8.2 \cdot 10^{-3} \times \left[ 1 - 0.8 \ln \left( \frac{m_N^2}{M_W^2} \right) + 0.5 \ln^2 \left( \frac{m_N^2}{M_W^2} \right) \right] B(\mu \rightarrow e\gamma) . \quad (12)$$

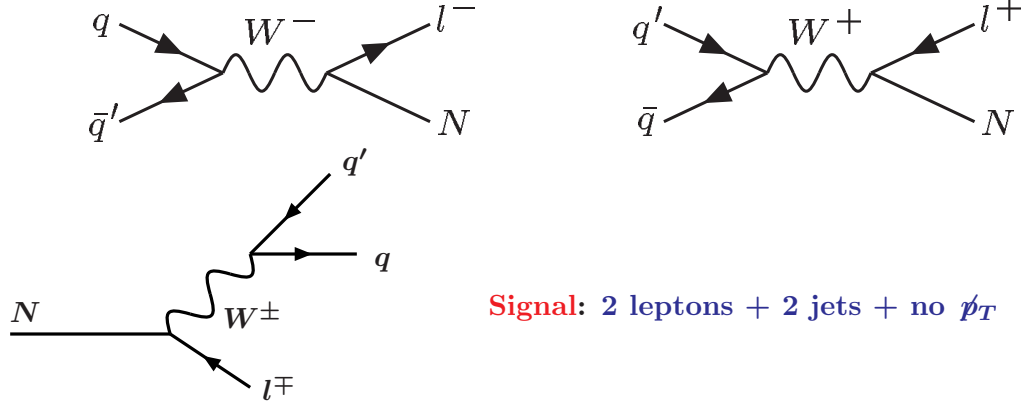
For example, for an R $\tau$ L model with  $m_N = 250$  GeV, (12) implies

$$B(\mu \rightarrow eee) \simeq 1.4 \cdot 10^{-2} \times B(\mu \rightarrow e\gamma) \simeq 1.4 \cdot 10^{-14} \quad (13)$$

This value is a factor  $\sim 70$  below the present experimental bound [21]:  $B_{\text{exp}}(\mu \rightarrow eee) < 1.0 \times 10^{-12}$ . It would be therefore encouraging, if higher sensitivity experiments could be designed to further probe this observable.

Finally, one of the most sensitive experiments to probe LFV is the coherent conversion of  $\mu \rightarrow e$  in nuclei. Specifically, for the  ${}^{48}_{22}\text{Ti}$  case,  $B_{\mu e}(26, 22)$  can be also directly linked to  $B(\mu \rightarrow e\gamma)$  through the relation [12]:

$$B_{\mu e}(26, 22) \simeq 0.1 \times \left[ 1 + 0.5 \ln \left( \frac{m_N^2}{M_W^2} \right) \right]^2 B(\mu \rightarrow e\gamma) . \quad (14)$$



**Figure 7.** Production and decay of heavy Majorana neutrinos at the LHC.

On the experimental side, the upper bound on the  $\mu \rightarrow e$  conversion rate in  $^{48}\text{Ti}$  [53] is

$$B_{\mu e}^{\text{exp}}(26, 22) < 4.3 \times 10^{-12}, \quad (15)$$

at the 90% CL. However, the proposed experiment PRISM/PRIME [54] will be sensitive to conversion rates of order  $10^{-18}$ .

Given (14) and  $m_N = 250$  GeV, the prediction for  $\mu \rightarrow e$  conversion in  $^{48}\text{Ti}$  is:

$$B_{\mu e}(26, 22) \simeq 0.46 \times B(\mu \rightarrow e\gamma) \sim 4.5 \times 10^{-13}. \quad (16)$$

The above prediction falls well within reach of the sensitivity proposed by the PRISM/PRIME collaborations.

As opposed to other scenarios of baryogenesis, leptogenesis models in general do not suffer from predicting a too large contribution to the electron electric dipole moment (EDM). This is because EDM effects first arise at two loops and are suppressed either by higher powers of small Yukawa couplings and/or by small factors, such as  $(m_{N_1} - m_{N_{2,3}})/m_N$  in RL models [8, 55]. Finally, other manifestations of heavy Majorana neutrino effects are LFV and universality-breaking effects in the  $Z$ -boson decays [49, 50] and non-unitarity effects in neutrino oscillations [56].

### 5.3. Heavy Majorana Neutrino Production at the LHC

Electroweak-scale RL models, such as the  $R\tau L$  model, predict heavy Majorana neutrinos with masses and couplings within the range that could be produced at the LHC. The dominant production channel is through the  $W^\pm$ -boson exchange graph shown in Fig. 7 [14, 15, 57, 58, 59]. The characteristic signature is the observation of like-sign dileptons with no missing transverse momentum [60].

According to the recent studies [57, 58], the LHC may be rather sensitive to the light-to-heavy neutrino mixing expression  $b^2 v^2 / (2m_N^2) \sim 10^{-3} - 10^{-4}$ , for  $m_N = 100 - 150$  GeV. Hence, a possible discovery of heavy Majorana neutrinos at the LHC will unravel the scale of lepton-flavour violation, thus putting the predictions of  $R\tau L$  models into direct test. Finally, low-scale seesaw models of the type considered here could also predict sizeable signatures of simultaneous lepton flavour and lepton number violation, such as  $pp \rightarrow e^+ \mu^+, e^- \mu^-, e^- \tau^- \dots$ , as well as resonantly enhanced  $\mathcal{O}(1)$  CP asymmetries [61]. Thus, a complete exploration of the phenomenological consequences of RL models would be interesting.

## 6. Conclusions

Leptogenesis provides an interesting mechanism for explaining our observable matter–antimatter asymmetric Universe, which may relate to the origin of the neutrino mass. In this brief review, we have paid particular attention to the recent developments on the field, such as the flavourdynamics of leptogenesis and resonant leptogenesis near the electroweak phase transition. In particular, flavourdynamics plays an important role in the predictions for the BAU and in the model-building of phenomenologically testable scenarios. It has been explicitly demonstrated how the observed matter–antimatter asymmetry in the Universe may be successfully explained by thermal resonant leptogenesis at the electroweak scale, independently of any non-zero primordial baryon asymmetry. In fact, any primordial baryon asymmetry will be erased and the actual value of the BAU will be set by the actual mechanism itself, as a consequence of the quasi-in-thermal equilibrium conditions that govern resonant leptogenesis.

Models of resonant leptogenesis may have rich collider phenomenology. Electroweak-mass heavy Majorana neutrinos may give rise to observable signatures of lepton number and lepton flavour violation at the LHC. There may also give rise to correlated predictions for low-energy observables, such as  $B(\mu \rightarrow e\gamma)$ , coherent  $\mu \rightarrow e$  conversion in nuclei,  $B(\mu \rightarrow eee)$  and the effective neutrino mass in the  $0\nu\beta\beta$  decay. Thus, an exciting era of new discoveries at the LHC and in other low-energy experiments may lie just ahead of us, revealing a greater picture that unifies particle physics and cosmology.

## Acknowledgements

I am grateful to the organisers of the conference DISCRETE'08 for their hospitality.

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