Flavour-Dependent Type II Leptogenesis

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Abstract

We reanalyse leptogenesis via the out-of-equilibrium decay of the lightest right-handed neutrino in type II seesaw scenarios, taking into account flavour-dependent effects. In the type II seesaw mechanism, in addition to the type I seesaw contribution, an additional direct mass term for the light neutrinos is present. We consider type II seesaw scenarios where this additional contribution arises from the vacuum expectation value of a Higgs triplet, and furthermore an effective model-independent approach. We investigate bounds on the flavour-specific decay asymmetries, on the mass of the lightest right-handed neutrino and on the reheat temperature of the early universe, and compare them to the corresponding bounds in the type I seesaw framework. We show that while flavour-dependent thermal type II leptogenesis becomes more efficient for larger mass scale of the light neutrinos, and the bounds become relaxed, the type I seesaw scenario for leptogenesis becomes more constrained. We also argue that in general, flavour-dependent effects cannot be ignored when dealing with leptogenesis in type II seesaw models.
1 Introduction

Leptogenesis [1] is one of the most attractive and minimal mechanisms for explaining the observed baryon asymmetry of the universe $n_B/n_\gamma \approx (6.0965 \pm 0.2055) \times 10^{-10}$ [2]. A lepton asymmetry is dynamically generated and then converted into a baryon asymmetry due to $(B + L)$-violating sphaleron interactions [3] which exist in the Standard Model (SM) and its minimal supersymmetric extension, the MSSM. Leptogenesis can be implemented within the type I seesaw scenario [4], consisting of the SM (MSSM) plus three right-handed Majorana neutrinos (and their superpartners) with a hierarchical spectrum. In thermal leptogenesis [5], the lightest of the right-handed neutrinos is produced by thermal scattering after inflation, and subsequently decays out-of-equilibrium in a lepton number and CP-violating way, thus satisfying Sakharov’s constraints [6].

In models with a left-right symmetric particle content like minimal left-right symmetric models, Pati-Salam models or Grand Unified Theories (GUTs) based on SO(10), the type I seesaw mechanism is typically generalised to a type II seesaw [7], where an additional direct mass term $m_{II}^{LL}$ for the light neutrinos is present. From a model independent perspective, the type II mass term can be considered as an additional contribution to the lowest dimensional effective neutrino mass operator. In most explicit models, the type II contribution stems from seesaw suppressed induced vevs of SU(2)$_L$-triplet Higgs fields. One motivation for considering the type II seesaw is that it allows to construct unified flavour models for partially degenerate neutrinos in an elegant way, e.g. via a type II upgrade [8], which is otherwise difficult to achieve in type I models.

For leptogenesis in type II seesaw scenarios with SU(2)$_L$-triplet Higgs fields, there are in general two possibilities to generate the baryon asymmetry: via decays of the lightest right-handed neutrinos or via decays of the SU(2)$_L$-triplets [9, 10, 11, 12]. In the first case, there are additional one-loop diagrams where virtual triplets are running in the loop [9, 13, 14, 15, 16]. In the following, we focus on this possibility, and assume hierarchical right-handed neutrino masses (and that the triplets are heavier than $\nu^L_1$). In this limit, to a good approximation the decay asymmetry depends mainly on the low energy neutrino mass matrix $m_{LL}^{II} = m_{II}^{LL} + m_{II}^{RL}$ and on the Yukawa couplings to the lightest right-handed neutrino and its mass [16]. It has been shown that type II leptogenesis imposes constraints on the seesaw parameters, which, in the flavour-independent approximation, differ substantially from the constraints in the type I case. For instance, the bound on the decay asymmetry increases with increasing neutrino mass scale [16], in contrast to the type I case where it decreases. As a consequence, the lower bound on the mass of the lightest right-handed neutrino from leptogenesis decreases for increasing neutrino mass scale [16]. One interesting application of type II leptogenesis is the possibility to improve consistency of classes of unified flavour models with respect to thermal leptogenesis [17]. Finally, since the type II contribution typically does not effect washout, there is no bound on the absolute neutrino mass scale from type II leptogenesis, as has been pointed out in [15]. For further applications and realisations of type II leptogenesis in specific models of fermion masses and mixings, see e.g. [18].

In recent years, the impact of flavour in thermal leptogenesis has merited increasing attention [19] - [38]. In fact, the one-flavour approximation is only rigorously correct
when the interactions mediated by the charged lepton Yukawa couplings are out of equilibrium. Below a given temperature (e.g. $\mathcal{O}(10^{12} \text{ GeV})$ in the SM and $(1 + \tan^2 \beta) \times \mathcal{O}(10^{12} \text{ GeV})$ in the MSSM), the tau Yukawa coupling comes into equilibrium (later followed by the couplings of the muon and electron). Flavour effects are then physical and become manifest, not only at the level of the generated CP asymmetries, but also regarding the washout processes that destroy the asymmetries created for each flavour. In the full computation, the asymmetries in each distinguishable flavour are differently washed out, and appear with distinct weights in the final baryon asymmetry.

Flavour-dependent leptogenesis in the type I seesaw scenario has recently been addressed in detail by several authors. In particular, flavour-dependent effects in leptogenesis have been studied, and shown to be relevant, in the two right-handed neutrino models [24] as well as in classes of neutrino mass models with three right-handed neutrinos [26]. The quantum oscillations/correlations of the asymmetries in lepton flavour space have been included in [22,32,33,35] and the treatment has been generalised to the MSSM [26,29]. Effects of reheating, and constraints on the seesaw parameters from upper bounds on the reheat temperature, have been investigated in [29]. Leptogenesis bounds on the reheat temperature [29] and on the mass of the lightest right-handed neutrino [20,36] have also been considered including flavour-dependent effects. Strong connections between the low-energy CP phases of the $U_{MNS}$ matrix and CP violation for flavour-dependent leptogenesis have been shown to emerge in certain classes of neutrino mass models [20] or under the hypothesis of no CP violation sources associated with the right-handed neutrino sector (real $R$) [25,27,28,31]. Possible effects regarding the decays of the heavier right-handed neutrinos for leptogenesis have been discussed in this context in [21,34], and flavour-dependent effects for resonant leptogenesis were addressed in [38]. Regarding the masses of the light neutrinos, assuming hierarchical right-handed neutrinos and considering experimentally allowed light neutrino masses (below about 0.4 eV), there is no longer a bound on the neutrino mass scale from thermal leptogenesis if flavour-dependent effects are included [24].

In view of the importance of flavour-dependent effects on leptogenesis in the type I seesaw case, it is pertinent to investigate their effects on type II leptogenesis. In this paper, we therefore reanalyse leptogenesis via the out-of-equilibrium decay of the lightest right-handed neutrino in type II seesaw scenarios, taking into account flavour-dependent effects. We investigate bounds on the decay asymmetries, on the mass of the lightest right-handed neutrino and on the reheat temperature of the early universe, and discuss how increasing the neutrino mass scale affects thermal leptogenesis in the type I and type II seesaw frameworks.

2 Type I and type II seesaw mechanisms

Motivated by left-right symmetric unified theories, we consider two generic possibilities for explaining the smallness of neutrino masses: via heavy SM (MSSM) singlet fermions (i.e. right-handed neutrinos) [4] and via heavy SU(2)$_L$-triplet Higgs fields [7]. In both cases, the effective dimension five operator for Majorana neutrino masses in the SM or
Figure 1: Generation of the dimension 5 neutrino mass operator in the type I seesaw mechanism.

the MSSM, respectively,

\[ \mathcal{L}_\kappa^{\text{SM}} = \frac{1}{4} \kappa_{gf} (\overline{L} C \cdot \phi) (L \cdot \phi) + \text{h.c.}, \quad (1a) \]

\[ \mathcal{L}_\kappa^{\text{MSSM}} = -\frac{1}{4} \kappa_{gf} (\hat{L}^g \cdot \hat{H}_u) (\hat{L}^f \cdot \hat{H}_u) \left| \theta \right\theta + \text{h.c.}, \quad (1b) \]

is generated from integrating out the heavy fields. This is illustrated in figures 1 and 2. In equation (1), the dots indicate the SU(2)\textsubscript{L}-invariant product, \((\hat{L}^f \cdot \hat{H}_u) = \hat{L}^f_a (i\tau^2)_{ab} (\hat{H}_u)_b\), with \(\tau_A \ (A \in \{1, 2, 3\})\) being the Pauli matrices. Superfields are marked by hats. After electroweak symmetry breaking, the operators of equation (1) lead to Majorana mass terms for the light neutrinos,

\[ \mathcal{L}_\nu = -\frac{1}{2} m_{\nu_{LL}}^\nu \overline{\nu}_L \nu^C_L, \text{ with } m_{\nu_{LL}}^\nu = -\frac{v_u^2}{2} (\kappa)^* . \quad (2) \]

In the type I seesaw mechanism, it is assumed that only the singlet (right-handed) neutrinos \(\nu_{Ri}\) contribute to the neutrino masses. With \(Y_\nu\) being the neutrino Yukawa matrix in left-right convention\(^1\) \(M_{RR}\) the mass matrix of the right-handed neutrinos and \(v_u = \langle \phi^0 \rangle = \langle H_u^0 \rangle\) the vacuum expectation value of the Higgs field which couples to the right-handed neutrinos, the effective mass matrix of the light neutrinos is given by the conventional type I seesaw formula

\[ m_{\nu_{LL}}^I = -v_u^2 Y_\nu M_{RR}^{-1} Y_\nu^T . \quad (3) \]

In the type II seesaw mechanism, the contributions to the neutrino mass matrix from both, right-handed neutrinos \(\nu_{Ri}\) and Higgs triplet(s) \(\Delta_L\), are considered. The additional contribution to the neutrino masses from \(\Delta_L\) can be understood in two ways: as another contribution to the effective neutrino mass operator in the low energy effective theory or, equivalently, as a direct mass term after the Higgs triplet obtains an induced small vev after electroweak symmetry breaking (c.f. figure 2). The neutrino mass matrix in the type II seesaw mechanism has the form

\[ m_{\nu_{LL}}^\nu = m_{\nu_{LL}}^I + m_{\nu_{LL}}^{II} = m_{\nu_{LL}}^I - v_u^2 Y_\nu M_{RR}^{-1} Y_\nu^T , \quad (4) \]

\(^1\)The neutrino Yukawa matrix corresponds to \(- (Y_\nu)_{fi} (L^f \cdot \phi) \nu_R^i\) in the Lagrangian of the SM and, analogously, to \((Y_\nu)_{fi} (L^f \cdot H_u) \nu^C_i\) in the superpotential of the MSSM (see [16] for further details).
where $m_{II}^{H}$ is the additional term from the Higgs triplet(s). In left-right symmetric unified theories, the generic size of both seesaw contributions $m_{II}^{L}$ and $m_{II}^{H}$ is $O(v_{u}^{2}/v_{B-L})$ where $v_{B-L}$ is the B-L breaking scale (i.e. the mass scale of the right-handed neutrinos and of the Higgs triplet(s)).

3 Baryogenesis via flavour-dependent leptogenesis

Flavour-dependent effects can have a strong impact in baryogenesis via thermal leptogenesis [19] - [38]. The effects are manifest not only in the flavour-dependent CP asymmetries, but also in the flavour-dependence of scattering processes in the thermal bath, which can destroy a previously produced asymmetry.

The relevance of the flavour-dependent effects depends on the temperatures at which thermal leptogenesis takes place, and thus on which interactions mediated by the charged lepton Yukawa couplings are in thermal equilibrium. For example, in the MSSM, for temperatures between circa $(1+\tan^{2}\beta)\times10^{5}\text{GeV}$ and $(1+\tan^{2}\beta)\times10^{9}\text{GeV}$, the $\mu$ and $\tau$ Yukawa couplings are in thermal equilibrium and all flavours in the Boltzmann equations are to be treated separately. For $\tan\beta = 30$, this applies for temperatures below about $10^{12}\text{GeV}$ and above $10^{8}\text{GeV}$, a temperature range which is of most interest for thermal leptogenesis in the MSSM. In the SM, in the temperature range between circa $10^{9}\text{GeV}$ and $10^{12}\text{GeV}$, only the $\tau$ Yukawa coupling is in equilibrium and is treated separately in the Boltzmann equations, whereas $\mu$ and $e$ flavours are indistinguishable. A discussion of the temperature regimes in the SM and MSSM, where flavour is important, can be found, e.g., in [26].

We now briefly review the estimation of the produced baryon asymmetry in flavour-dependent leptogenesis\textsuperscript{2}. For definiteness, we focus on the temperature range where all flavours are to be treated separately. In the following discussion of thermal type II leptogenesis, we will assume that the mass $M_{\Delta L}$ of the triplet(s) is much larger than $M_{R1}$. In this limit, the flavour-dependent efficiencies calculated in the type I seesaw scenario can also be used in the type II framework. The out-of-equilibrium decays of the heavy right-handed (s)neutrinos $\nu_{R}^{1}$ and $\tilde{\nu}_{R}^{1}$ give rise to flavour-dependent asymmetries in the (s)lepton sector, which are then partly transformed via sphaleron conversion into

\textsuperscript{2}For a discussion of approximations which typically enter these estimates, and which also apply to our discussion, see e.g. section 3.1.3 in [29].
a baryon asymmetry $Y_B$. The final baryon asymmetry can be calculated as

$$Y_B^{\text{SM}} = \frac{12}{37} \sum_f Y_{\Delta f}^{\text{SM}},$$

$$Y_B^{\text{MSSM}} = \frac{10}{31} \sum_f \hat{Y}_{\Delta f}^{\text{MSSM}},$$

where $\hat{Y}_{\Delta f} \equiv Y_B/3 - Y_L$ are the total (particle and particle) $B/3 - L_f$ asymmetries, with $Y_{L_f}$ the lepton number densities in the flavour $f = e, \mu, \tau$. The asymmetries $Y_{\Delta f}^{\text{MSSM}}$ and $Y_{\Delta f}^{\text{SM}}$, which are conserved by sphalerons and by the other SM (MSSM) interactions, are then usually calculated by solving a set of coupled Boltzmann equations, describing the evolution of the number densities as a function of temperature.

It is convenient to parameterise the produced asymmetries in terms of flavour-specific efficiency factors $\eta_f$ and decay asymmetries $\varepsilon_{1,f}$ as

$$Y_{\Delta f}^{\text{SM}} = \eta_{f}^{\text{SM}} \varepsilon_{1,f} Y_{\nu_R}^{\text{eq}},$$

$$\hat{Y}_{\Delta f}^{\text{MSSM}} = \eta_{f}^{\text{MSSM}} \left[ \frac{1}{2}(\varepsilon_{1,f} + \varepsilon_{1,f}) Y_{\nu_R}^{\text{eq}} + \frac{1}{2}(\varepsilon_{1,f} - \varepsilon_{1,f}) Y_{\tilde{\nu}_R}^{\text{eq}} \right].$$

$Y_{\nu_R}^{\text{eq}}$ and $Y_{\tilde{\nu}_R}^{\text{eq}}$ are the number densities of the neutrino and sneutrino for $T \gg M_1$ if they were in thermal equilibrium, normalised with respect to the entropy density. In the Boltzmann approximation, they are given by $Y_{\nu_R}^{\text{eq}} \approx Y_{\tilde{\nu}_R}^{\text{eq}} \approx 45/(\pi^4 g\ast)$. $g\ast$ is the effective number of degrees of freedom, which amounts 106.75 in the SM and 228.75 in the MSSM.

$\varepsilon_{1,f}, \varepsilon_{1,f}, \varepsilon_{1,f}, \varepsilon_{1,f}$ are the decay asymmetries for the decay of neutrino into Higgs and lepton, neutrino into Higgsino and slepton, sneutrino into Higgsino and lepton, and sneutrino into Higgs and slepton, respectively, defined by

$$\varepsilon_{1,f} = \frac{\Gamma_{\nu_R^L L_f} - \Gamma_{\nu_R^L L_f}}{\sum_f (\Gamma_{\nu_R^L L_f} + \Gamma_{\nu_R^L L_f})}, \quad \varepsilon_{1,f} = \frac{\Gamma_{\nu_R^L L_f} - \Gamma_{\nu_R^L L_f}}{\sum_f (\Gamma_{\nu_R^L L_f} + \Gamma_{\nu_R^L L_f})},$$

$$\varepsilon_{1,f} = \frac{\Gamma_{\tilde{\nu}_R^L L_f} - \Gamma_{\tilde{\nu}_R^L L_f}}{\sum_f (\Gamma_{\tilde{\nu}_R^L L_f} + \Gamma_{\tilde{\nu}_R^L L_f})}, \quad \varepsilon_{1,f} = \frac{\Gamma_{\tilde{\nu}_R^L L_f} - \Gamma_{\tilde{\nu}_R^L L_f}}{\sum_f (\Gamma_{\tilde{\nu}_R^L L_f} + \Gamma_{\tilde{\nu}_R^L L_f})}.$$
Figure 3: Flavour-dependent efficiency factor $\eta(A_{f f} K_f, K)$ in the MSSM as a function of $A_{f f} K_f$, for fixed values of $K/|A_{f f} K_f| = 2, 5$ and $100$, obtained from solving the flavour-dependent Boltzmann equations in the MSSM with zero initial abundance of right-handed (s)neutrinos (figure from [26]). $A$ is a matrix which appears in the Boltzmann equations (see [19,24] for $A$ in the SM and [26] for the MSSM case), and which has diagonal elements $|A_{f f}|$ of $O(1)$. The small off-diagonal entries of $A$ have been neglected, which is a good approximation in most cases. In general, however, they have to be included. More relevant than the differences in the flavour-dependent efficiency factors for different $K/|A_{f f} K_f|$ is that the total baryon asymmetry is the sum of each individual lepton asymmetries, which is weighted by the corresponding efficiency factors.

calculated in the type I seesaw scenario can also be used in the type II framework. In the definition of the efficiency factor, the equilibrium number densities serve as a normalization: A thermal population $\nu_{R1}$ (and $\tilde{\nu}_{R1}$) decaying completely out of equilibrium (without washout effects) would lead to $\eta_f = 1$.

The efficiency factors can be computed by means of the flavour-dependent Boltzmann equations, which can be found for the SM in [19,22,23,24] and for the MSSM in [26,29]. In general, the flavour-dependent efficiencies depend strongly on the washout parameters $\tilde{m}_{1,f}$ for each flavour, and on the total washout parameter $\tilde{m}_1$, which are defined as

$$\tilde{m}_{1,f} = \frac{v^2_u |(Y_\nu)_f|^2}{M_{R1}}, \quad \tilde{m}_1 = \sum_f \tilde{m}_{1,f}. \quad (10)$$

Alternatively, one may use the quantities $K_f, K$, which are related to $\tilde{m}_{1,f}, \tilde{m}_1$ by

$$K_f = \frac{\tilde{m}_{1,f}}{m^*}, \quad K = \sum_f K_f, \quad (11)$$

with $m^*_{\text{SM}} \approx 1.08 \times 10^{-3}$ eV and $m^*_{\text{MSSM}} \approx \sin^2(\beta) \times 1.58 \times 10^{-3}$ eV. Figure 3 shows the flavour-specific efficiency factor $\eta_f$ in the MSSM. Maximal efficiency for a specific flavour corresponds to $K_f \approx 1$ ($\tilde{m}_{1,f} \approx m^*$).

The most relevant difference between the flavour-independent approximation and the correct flavour-dependent treatment is the fact that in the latter, the total baryon asymmetry is the sum of each individual lepton asymmetries, which is weighted by the
corresponding efficiency factor. Therefore, upon summing over the lepton asymmetries, the total baryon number is generically not proportional to the sum over the CP asymmetries, $\varepsilon_1 = \sum_f \varepsilon_{1,f}$, as in the flavour-independent approximation where the lepton flavour is neglected in the Boltzmann equations. In other words, in the flavour-independent approximation the total baryon asymmetry is a function of $(\sum_f \varepsilon_{1,f})$ × $\eta^{\text{ind}} (\sum_g K_g)$. In the correct flavour treatment the baryon asymmetry is (approximately) a function of $\sum_f \varepsilon_{1,f} (\eta) (A_{\text{ff}} K_f, K)$. From this, it is already clear that flavour-dependent effects can have important consequences also in type II leptogenesis.

The most important quantities for computing the produced baryon asymmetry are thus the decay asymmetries $\varepsilon_{1,f}$ and the efficiency factors $\eta_f$ (which depend mainly on $\tilde{m}_{1,f}$ and $\tilde{m}_1$ (or $K_f$ and $K$)). While the efficiency factors can be computed similarly to the type I seesaw case, important differences between leptogenesis in type I and type II seesaw scenarios arise concerning the decay asymmetries as well as concerning the connection between leptogenesis and seesaw parameters.

4 Decay asymmetries

4.1 Right-handed neutrinos plus triplets

Regarding the decay asymmetry in the type II seesaw mechanism, where the direct mass term for the neutrinos stems from the induced vev of a Higgs triplet, there are new contributions from 1-loop diagrams where virtual SU(2)$_L$-triplet scalar fields (or their superpartners) are exchanged in the loop. The relevant diagrams for the decay $\nu_R^1 \rightarrow L_a^\ell H_{ab}$ in the limit $M_1 \ll M_{R2}, M_{R1}, M_\Delta$ are shown in figure 4. Compared to the type I seesaw framework, the new contributions are the diagrams (c) and (f). The calculation of the corresponding decay asymmetries for each lepton flavour yields

$$
\varepsilon_{1,f}^{(a)} = \frac{1}{8\pi} \sum_{j \neq f} \frac{\text{Im} [(Y^\ell)^{1f} (Y^\nu)^{1j} (Y_T)^{jf}]}{v_{\nu_j} (Y^\nu)^{1j}} \sqrt{x_j} \left[ 1 - (1 + x_j) \ln \left( \frac{x_j + 1}{x_j} \right) \right],
$$

$$
\varepsilon_{1,f}^{(b)} = \frac{1}{8\pi} \sum_{j \neq f} \frac{\text{Im} [(Y^\ell)^{1f} (Y^\nu)^{1j} (Y_T)^{jf}]}{v_{\nu_j} (Y^\nu)^{1j}} \sqrt{x_j} \left[ \frac{1}{1 - x_j} \right],
$$

$$
\varepsilon_{1,f}^{(c)} = -\frac{3}{8\pi} M_{R1} x_{\nu_j} \text{Im} [(Y^\ell)^{1g} (Y^\nu)^{1j} (m_{LL}^H)_{gf}] y \left[ -1 + y \ln \left( \frac{y + 1}{y} \right) \right],
$$

$$
\varepsilon_{1,f}^{(d)} = \frac{1}{8\pi} \sum_{j \neq f} \frac{\text{Im} [(Y^\ell)^{1f} (Y^\nu)^{1j} (Y_T)^{jf}]}{v_{\nu_j} (Y^\nu)^{1j}} \sqrt{x_j} \left[ -1 + x_j \ln \left( \frac{x_j + 1}{x_j} \right) \right],
$$

$$
\varepsilon_{1,f}^{(e)} = \frac{1}{8\pi} \sum_{j \neq f} \frac{\text{Im} [(Y^\ell)^{1f} (Y^\nu)^{1j} (Y_T)^{jf}]}{v_{\nu_j} (Y^\nu)^{1j}} \sqrt{x_j} \left[ \frac{1}{1 - x_j} \right],
$$

$$
\varepsilon_{1,f}^{(f)} = -\frac{3}{8\pi} M_{R1} x_{\nu_j} \text{Im} [(Y^\ell)^{1g} (Y^\nu)^{1j} (m_{LL}^H)_{gf}] y \left[ 1 - (1 + y) \ln \left( \frac{y + 1}{y} \right) \right].
$$
Figure 4: Loop diagrams in the MSSM which contribute to the decay $\nu_R^1 \rightarrow L^f H_{ub}$ for the case of a type II seesaw mechanism where the direct mass term for the neutrinos stems from the induced vev of a Higgs triplet. In diagram (f), $\Delta_1$ and $\Delta_2$ are the mass eigenstates corresponding to the superpartners of the SU(2)$_L$-triplet scalar fields $\Delta$ and $\bar{\Delta}$. The SM diagrams are the ones where no superpartners (marked by a tilde) are involved and where $H_u$ is renamed to the SM Higgs $\phi$.

where $y := M_\Delta^2/M_{R1}^2$ and $x_j := M_{Rj}^2/M_{R1}^2$ for $j \neq 1$ and where we assume hierarchical right-handed neutrino masses and $M_\Delta \gg M_{R1}$.

The MSSM results for the type II contributions have been derived in [16]. In the SM, the results in [16] correct the previous result of [15] by a factor of $-3/2$. In equation (12) they have been generalised to the flavour-dependent case. The results for the contributions to the decay asymmetries from the triplet in the SM and from the triplet superfield in the MSSM are

$$\epsilon_{1,f}^{\text{SM,II}} = \epsilon_{1,f}^{(c)} ,$$
$$\epsilon_{1,f}^{\text{MSSM,II}} = \epsilon_{1,f}^{(c)} + \epsilon_{1,f}^{(f)} .$$

In the MSSM, we furthermore obtain

$$\epsilon_{1,f}^{\text{MSSM,II}} = \epsilon_{1,f}^{\text{MSSM,II}} = \epsilon_{1,f}^{\text{MSSM,II}} = \epsilon_{1,f}^{\text{MSSM,II}} .$$

The results corresponding to the diagrams (a), (b), (d) and (e) which contribute to $\epsilon_1$ in the type I seesaw in the SM and in the MSSM, have been presented first in [39]. The results for the type I contribution to the decay asymmetries in the SM and in the

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Integrating out the heavy particles $\nu_R^2, \nu_R^3, \Delta$ (and their superpartners) in figure 4 leads to an effective approach involving the dimension 5 neutrino mass operator (c.f. figures 1, 2 and 5), as will be discussed in section 4.2. We note that there are additional diagrams not shown in figure 4 (since they are generically suppressed for $M_1 \ll M_{R2, M_{R2}, M_\Delta}$) which are related to the dimension 6 operator containing two lepton doublets, two Higgs doublets and a derivative.

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4Integrating out the heavy particles $\nu_R^2, \nu_R^3, \Delta$ (and their superpartners) in figure 4 leads to an effective approach involving the dimension 5 neutrino mass operator (c.f. figures 1, 2 and 5), as will be discussed in section 4.2. We note that there are additional diagrams not shown in figure 4 (since they are generically suppressed for $M_1 \ll M_{R2, M_{R2}, M_\Delta}$) which are related to the dimension 6 operator containing two lepton doublets, two Higgs doublets and a derivative.
MSSM are

\[ \varepsilon_{\text{SM},1,f}^{\text{I}} = \varepsilon_{1,f}^{(a)} + \varepsilon_{1,f}^{(b)} , \]

\[ \varepsilon_{\text{MSSM},1,f}^{\text{I}} = \varepsilon_{1,f}^{(a)} + \varepsilon_{1,f}^{(b)} + \varepsilon_{1,f}^{(c)} . \]

(15a)

(15b)

Again, in the MSSM, the remaining decay asymmetries are equal to \( \varepsilon_{1,f}^{\text{MSSM}} \):

\[ \varepsilon_{1,f}^{\text{MSSM}} = \varepsilon_{\text{MSSM},1,f}^{\text{I}} = \varepsilon_{\text{MSSM},1,f}^{\text{II}} = \varepsilon_{1,f}^{\text{MSSM}} . \]

(16)

Finally, the total decay asymmetries from the decay of \( \nu_{1}^{R} \) in the type II seesaw, where the direct mass term for the neutrinos stems from the induced vev of a Higgs triplet, are given by

\[ \varepsilon_{1,f}^{\text{SM}} = \varepsilon_{1,f}^{\text{SM},1} + \varepsilon_{1,f}^{\text{SM},2} , \]

\[ \varepsilon_{1,f}^{\text{MSSM}} = \varepsilon_{1,f}^{\text{MSSM},1} + \varepsilon_{1,f}^{\text{MSSM},2} . \]

(17)

(18)

It is interesting to note that the type I results can be brought to a form which contains the neutrino mass matrix using

\[ \frac{\sum_{j \neq 1} \text{Im} \left[ (Y_{1j}^{*})_{ij} (Y_{1j}^{*})_{ij} \right]}{8 \pi (Y_{1j}^{*})_{11} \sqrt{x_{j}}} = - \frac{M_{R1}}{v_{u}^{2}} \frac{\sum_{g} \text{Im} \left[ (Y_{1g}^{*})_{ij} (Y_{1g}^{*})_{g1} (m_{1}^{I} + m_{2}^{I})_{fg} \right]}{8 \pi (Y_{1j}^{*})_{11}} . \]

(19)

In the limit \( y \gg 1 \) and \( x_{j} \gg 1 \) for all \( j \neq 1 \), which corresponds to a large gap between the mass \( M_{R1} \) and the masses \( M_{R2} \), \( M_{R3} \) and \( M_{\Delta} \), we obtain the simple results for the flavour-specific decay asymmetries \( \varepsilon_{1,f}^{\text{SM}} \) and \( \varepsilon_{1,f}^{\text{MSSM}} \) [16]

\[ \varepsilon_{1,f}^{\text{SM}} = \frac{3}{16 \pi} \frac{M_{R1}}{v_{u}^{2}} \frac{\sum_{g} \text{Im} \left[ (Y_{1g}^{*})_{ij} (Y_{1g}^{*})_{g1} (m_{1}^{I} + m_{2}^{I})_{fg} \right]}{(Y_{1j}^{*})_{11}} , \]

(20a)

\[ \varepsilon_{1,f}^{\text{MSSM}} = \frac{3}{8 \pi} \frac{M_{R1}}{v_{u}^{2}} \frac{\sum_{g} \text{Im} \left[ (Y_{1g}^{*})_{ij} (Y_{1g}^{*})_{g1} (m_{1}^{I} + m_{2}^{I})_{fg} \right]}{(Y_{1j}^{*})_{11}} . \]

(20b)

In the presence of such a mass gap, the calculation can also be performed in an effective approach after integrating out the two heavy right-handed neutrinos and the heavy triplet, as we now discuss.

### 4.2 Effective approach to leptogenesis

Let us now explicitly use the assumption that the lepton asymmetry is generated via the decay of the lightest right-handed neutrino and that all other additional particles, in particular the ones which generate the type II contribution, are much heavier than \( M_{R1} \). Furthermore, we assume that we can neglect their population in the early universe, e.g. that their masses are much larger than the reheat temperature \( T_{RH} \) and that they are
Figure 5: Loop diagrams contributing to the decay asymmetry via the decay $\nu^L_1 \to L^f H_{ub}$ in the MSSM with a (lightest) right-handed neutrino $\nu^R_1$ and a neutrino mass matrix determined by $\kappa^\prime$ \cite{16}. Further contributions to the generated baryon asymmetry stem from the decay of $\nu^R_1$ into slepton and Higgsino and from the decays of the sneutrino $\tilde{\nu}^L_1$. With $H_u$ renamed to the SM Higgs, the first diagram contributes in the extended SM.

not produced non-thermally in a large amount. Under these assumptions we can apply an effective approach to leptogenesis, which is independent of the mechanism which generates the additional (type II) contribution to the neutrino mass matrix \cite{16}.

For this minimal effective approach, it is convenient to isolate the type I contribution from the lightest right-handed neutrino as follows:

$$m_{LL}^{\nu} f g = \frac{-\nu^2_u}{2} \left[ 2(Y^\nu)_{f1} M_{R1}^{-1} (Y^T_{\nu})_{1g} + (\kappa'^*)_{fg} \right].$$  \hspace{1cm} (21)

$\kappa'$ includes type I contributions from the heavier right-handed neutrinos, plus any additional (type II) contributions from heavier particles. Examples for realisations of the neutrino mass operator can be found, e.g., in \cite{40}.

At $M_{R1}$, the minimal effective field theory extension of the SM (MSSM) for leptogenesis includes the effective neutrino mass operator $\kappa'$ plus one right-handed neutrino $\nu^R_1$ with mass $M_{R1}$ and Yukawa couplings $(Y^\nu)_{f1}$ to the lepton doublets $L^f$, defined as $-(Y^\nu)_{f1}(L^f \cdot \phi) \nu^R_1$ in the Lagrangian of the SM and, analogously, as $(Y^\nu)_{f1}(\tilde{L}^f \cdot \tilde{H}_u) \nu^{C1}$ in the superpotential of the MSSM.

The contributions to the decay asymmetries in the effective approach stem from the interference of the diagram(s) for the tree-level decay of $\nu^R_1$ (and $\tilde{\nu}^R_1$) with the loop diagrams containing the effective operator, shown in figure 5. In the SM, we obtain the simple result \cite{16} for the flavour-specific effective decay asymmetries (corresponding to diagram (a) of figure 5)

$$\epsilon_{1,f}^{\text{SM}} = \frac{3}{16\pi} \frac{M_{R1} \sum_g \text{Im} [(Y^\nu^*)_{f1}(Y^\nu^*)_{g1}(m_{LL}^{\nu})_{fg}]}{(Y^\nu^* Y^\nu)_{11}}.$$  \hspace{1cm} (22)

For the supersymmetric case, diagram (a) and diagram (b) contribute to $\epsilon_{1,f}^{\text{MSSM}}$ and we obtain \cite{16}:

$$\epsilon_{1,f}^{\text{MSSM}} = \frac{3}{8\pi} \frac{M_{R1} \sum_g \text{Im} [(Y^\nu^*)_{f1}(Y^\nu^*)_{g1}(m_{LL}^{\nu})_{fg}]}{(Y^\nu Y^\nu)_{11}}.$$  \hspace{1cm} (23)
Explicit calculation furthermore yields

\[ \varepsilon_{1,f}^{\text{MSSM}} = \varepsilon_{1,f}^{\text{SM}} = \varepsilon_{1,f}^{\text{MSSM}}. \]  

(24)

The results are independent of the details of the realisation of the neutrino mass operator \( \kappa' \). Note that, since the diagrams where the lightest right-handed neutrino runs in the loop do not contribute to leptogenesis, we have written \( m_{\nu_{\ell}} = -v_\ell^2(\kappa')^*/2 \) instead of \( m_{\nu_{\ell}} := -v_\ell^2(\kappa')^*/2 \) in the formulae in equations (22) - (23). The decay asymmetries are directly related to the neutrino mass matrix \( m_{\nu_{LL}}^{\nu} \).

For neutrino masses via the type I seesaw mechanism, the results are in agreement with the known results \[39\], in the limit \( M_{R2}, M_{R3} \gg M_{R1} \). The results obtained in the effective approach are also in agreement with our full theory calculation in the type II scenarios with SU(2)\(_L\)-triplets in equation (12) \[16\], in the limit \( M_\Delta \gg M_{R1} \).

5 Type II bounds on decay asymmetries and on \( M_{R1} \)

In the limit \( M_{R2}, M_{R3}, M_\Delta \gg M_{R1} \) (or alternatively in the effective approach), upper bounds for the total decay asymmetries in type II leptogenesis, i.e. for the sums \( |\varepsilon_{1,f}^{\text{SM}}| = |\sum_f \varepsilon_{1,f}^{\text{SM}}| \) and \( |\varepsilon_{1,f}^{\text{MSSM}}| = |\sum_f \varepsilon_{1,f}^{\text{MSSM}}| \), have been derived in \[16\]. For the flavour-specific decay asymmetries \( \varepsilon_{1,f}^{\text{SM}} \) and \( \varepsilon_{1,f}^{\text{MSSM}} \), the bounds can readily be obtained as

\[ |\varepsilon_{1,f}^{\text{SM}}| \leq \frac{3}{16\pi} \frac{M_{R1}}{v_\ell^2} m_\nu^{\nu_{\max}} \]  

\[ |\varepsilon_{1,f}^{\text{MSSM}}| \leq \frac{3}{8\pi} \frac{M_{R1}}{v_\ell^2} m_\nu^{\nu_{\max}}. \]

(25)

They are thus identical to the bounds for the total asymmetries. In particular, they also increase with increasing mass scale of the light neutrinos. Note that, compared to the low energy value, the neutrino masses at the scale \( M_{R1} \) are enlarged by renormalization group running by \( \approx +20\% \) in the MSSM and \( \approx +30\% \) in the SM, which raises the bounds on the decay asymmetries by the same values (see e.g. figure 4 of \[41\]).

A situation where an almost maximal baryon asymmetry is generated by thermal leptogenesis can be realised, for example, if the total decay asymmetry nearly saturates its upper bound and if, in addition, the washout parameters \( \tilde{m}_{1,f} \) for all three flavours approximately take its optimal value. Classes of type II seesaw models, where this can be accommodated, have been considered in \[8,12,17\]. In these so-called “type-II-upgraded” seesaw models, the type II contribution to the neutrino mass matrix is proportional to the unit matrix (enforced e.g. by an SO(3) flavour symmetry or by one of its non-Abelian discrete subgroups). From equation (20), one can readily see that if the type II contribution \( (x 1) \) dominates the neutrino mass matrix \( m_{\nu_{LL}}^{\nu} \), and if \( (Y_{\nu})_{f1} \) are approximately equal for all flavours \( f = 1, 2, 3 \) and chosen such that the resulting \( \tilde{m}_{1,f} \) are approximately equal to \( m^* \), we have realised \( \eta_f \approx \eta_{\max} \) for all flavours and simultaneously nearly saturated the bound for the total decay asymmetry.

\[ ^5 \text{We further note that the bound for one of the flavour-specific decay asymmetries can be nearly saturated in this scenario if, for instance, } (Y_{\nu})_{21} \approx (Y_{\nu})_{31} \approx 0. \]
Figure 6: Bound on the decay asymmetry $\varepsilon_{1,f}$ in type II leptogenesis (solid blue line) and type I leptogenesis (dotted red line) as a function of the mass of the lightest neutrino $m_{\nu_{\text{min}}} := \min(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$ in type I and type II seesaw scenarios (see also [29]). The washout parameter $|A_{ff}|\tilde{m}_{1,f}$ is fixed to $m^*$ (close to optimal), and the asymmetry is normalised to $\varepsilon_{1,\text{max},0} = 3M_{R1}(\Delta m_{31}^2)_{1/2}/(16\pi v_u^2)$, where $\Delta m_{31}^2 \approx 2.5 \times 10^{-3}$ eV$^2$ is the atmospheric neutrino mass squared difference. We have considered the MSSM with $\tan \beta = 30$ as an explicit example.

Assuming a maximal efficiency factor $\eta_{\text{max}}$ for all flavours in a given scenario, and taking an upper bound for the masses of the light neutrinos $m_{\nu_{\text{max}}}$ as well as the observed value $n_B/n_\gamma \approx (6.0965 \pm 0.2055) \times 10^{-10}$ [2] for the baryon asymmetry, equation (25) can be transformed into lower type II bounds for the mass of the lightest right-handed neutrino [16]:

$$M_{R1}^{\text{SM}} \geq \frac{16\pi}{3} \frac{v_u^2}{m_{\nu_{\text{max}}}} \frac{n_B/n_\gamma}{0.99 \cdot 10^{-2} \eta_{\text{max}}}, \quad M_{R1}^{\text{MSSM}} \geq \frac{8\pi}{3} \frac{v_u^2}{m_{\nu_{\text{max}}}} \frac{n_B/n_\gamma}{0.92 \cdot 10^{-2} \eta_{\text{max}}}.$$  \hspace{2cm} (26)

The bound on $M_{R1}$ is lower for a larger neutrino mass scale.

The situation in the type II framework differs from the type I seesaw case: In the latter, the flavour-specific decay asymmetries are constrained by [24]

$$|\varepsilon_{1,f}^{\text{SM}}| \leq \frac{3}{16\pi} \frac{M_{R1}}{v_u^2} m_{\nu_{\text{max}}} \left(\frac{\tilde{m}_{1,f}}{\tilde{m}_1}\right)^{1/2}, \quad |\varepsilon_{1,f}^{\text{MSSM}}| \leq \frac{3}{8\pi} \frac{M_{R1}}{v_u^2} m_{\nu_{\text{max}}} \left(\frac{\tilde{m}_{1,f}}{\tilde{m}_1}\right)^{1/2}. \hspace{2cm} (27)$$

Note that compared to the type II bounds, there is an extra factor of $(\tilde{m}_{1,f}/\tilde{m}_1)^{1/2}$, which depends on the washout parameters. As we shall now discuss, this factor implies that it is not possible to have a maximal decay asymmetry $\varepsilon_{1,f}$ and an optimal washout parameter $\tilde{m}_{1,f}$ simultaneously. Let us recall first that in the type I seesaw, in contrast to the type II case, the flavour-independent washout parameter has the lower bound [43]

$$\tilde{m}_1 \geq m_{\nu_{\text{min}}}, \hspace{2cm} (28)$$

with $m_{\nu_{\text{min}}} = \min(m_{\nu_1}, m_{\nu_2}, m_{\nu_3})$. On the contrary, in the type I and type II seesaw, the flavour-dependent washout parameters $\tilde{m}_{1,f}$ are generically not constrained. Note that in
Figure 7: Lower bound on $M_{R1}$ in type II leptogenesis (solid blue line) and type I leptogenesis (dotted red line) as a function of the mass of the lightest neutrino $m_{\nu_{\text{min}}}^\nu := \min (m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$. For definiteness, the MSSM with $\tan \beta = 30$ has been considered as an example.

The flavour-independent approximation, Eq. (28), leads to a dramatically more restrictive bound on $\varepsilon_1 = \sum_f \varepsilon_{1,f}$ \cite{44} for quasi-degenerate light neutrino masses, and finally even to a bound on the neutrino mass scale \cite{43}. This can be understood from the fact that for $\tilde{m}_1 \gg m^*$ in the flavour-independent approximation, washout effects strongly reduce the efficiency of thermal leptogenesis. Similarly, in the flavour-dependent treatment, $\tilde{m}_{1,f} \gg m^*$ would lead to a strongly reduced efficiency for this specific flavour. This strong washout for quasi-degenerate light neutrinos can be avoided in flavour-dependent type I leptogenesis, and $\tilde{m}_{1,f} \approx m^*$ can realise a nearly optimal scenario regarding washout (c.f. figure 3). However, we see from equation (27) that the decay asymmetries in this case are reduced by a factor of $(m^*/m_{\nu_{\text{min}}}^\nu)^{1/2}$ when compared to the optimal value, leading to a reduced baryon asymmetry. On the other hand, realizing nearly optimal $\varepsilon_{1,f}$ requires $\tilde{m}_{1,f} \approx \tilde{m}_1 \geq m_{\nu_{\text{min}}}^\nu$, leading to large washout effects for quasi-degenerate light neutrinos and even to a more strongly suppressed generation of baryon asymmetry (c.f. figure 3). As a consequence, increasing the neutrino mass scale increases the lower bound on $M_{R1}$ (also in the presence of flavour-dependent effects), in contrast to the type II seesaw case.

Comparing the type II and type I seesaw cases, in the latter the baryon asymmetry is suppressed for quasi-degenerate light neutrino masses either by a factor $(m^*/m_{\nu_{\text{min}}}^\nu)^{1/2}$ in the decay asymmetries or by a non-optimal washout parameter much larger than $m^*$ (or $K_f \gg 1$, c.f. figure 3). The bounds on the decay asymmetries in type I and type II leptogenesis are compared in figure 6 where $\tilde{m}_{1,f}$ has been fixed to $m^*$, close to its optimal value. From figure 5 we see that in the type I case the maximal baryon asymmetry is obtained for hierarchical neutrino masses, whereas in the type II case, increasing the neutrino mass scale increases the produced baryon asymmetry and therefore allows to relax the bound on $M_{R1}$, as shown in figure 7. In addition, for the same reason, increasing the neutrino mass scale also relaxes the lower bound on the reheat temperature $T_{R_H}$ from the requirement of successful type II leptogenesis. Including reheating in the flavour-dependent Boltzmann equations as in Ref. 29 (for the flavour-independent
case, see [45]), we obtain the $m_{\nu_{\text{min}}}^\nu$-dependent lower bounds on $T_{RH}$ in type I and type II scenarios shown in figure 8. While the bound decreases in type II leptogenesis by about an order of magnitude when the neutrino mass scale increases to 0.4 eV, it increases in the type I seesaw case. In the presence of upper bounds on $T_{RH}$, this can lead to constraints on the neutrino mass scale, i.e. on $m_{\nu_{\text{min}}}^\nu = \min (m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$. For instance, with an upper bound $T_{RH} \leq 5 \times 10^9$ GeV, values of $m_{\nu_{\text{min}}}^\nu$ in the approximate range [0.01 eV, 0.32 eV] would be incompatible with leptogenesis in the type I seesaw framework (c.f. figure 8).

6 Summary, discussion and conclusions

We have analysed flavour-dependent leptogenesis via the out-of-equilibrium decay of the lightest right-handed neutrino in type II seesaw scenarios, where, in addition to the type I seesaw, an additional direct mass term for the light neutrinos is present. We have considered type II seesaw scenarios where this additional contribution stems from the vacuum expectation value of a Higgs triplet, and furthermore an effective approach, which is independent of the mechanism which generates the additional (type II) contribution to neutrino masses. We have taken into account flavour-dependent effects, which are relevant if thermal leptogenesis takes place at temperatures below circa $10^{12}$ GeV in the SM and below circa $(1 + \tan^2 \beta) \times 10^{12}$ GeV in the MSSM. As in type I leptogenesis, in the flavour-dependent regime the decays of the right-handed (s)neutrinos generate asymmetries in in each distinguishable flavour (proportional to the flavour-specific decay asymmetries $\varepsilon_{1,f}$), which are differently washed out by scattering processes in the thermal bath, and thus appear with distinct weights (efficiency factors $\eta_f$) in the final baryon asymmetry.

The most important quantities for computing the produced baryon asymmetry are
the decay asymmetries $\varepsilon_{1,f}$ and the efficiency factors $\eta_f$ (which mainly depend on washout parameters $\tilde{m}_{1,f}$ and $\tilde{m}_1 = \sum_f \tilde{m}_{1,f}$). With respect to the flavour-specific efficiency factors $\eta_f$, in the limit that the mass $M_{\Delta L}$ of the triplet is much larger than $M_{R1}$ (and $M_{R1} \ll M_{R2}, M_{R3}$), they can be estimated from the same Boltzmann equations as in the type I seesaw framework. Regarding the decay asymmetries $\varepsilon_{1,f}$, in the type II seesaw case there are additional contributions where virtual Higgs triplets (and their superpartners) run in the 1-loop diagrams. Here, we have generalised the results of [16] to the flavour-dependent case. The most important effects of flavour in leptogenesis are a consequence of the fact that in the flavour-independent approximation the total baryon asymmetry is a function of $\left( \sum_f \tilde{m}_{1,f} \right) \times \eta^{\text{ind}} \left( \sum_g \tilde{m}_{1,g} \right)$, whereas in the correct flavour-dependent treatment the baryon asymmetry is (approximately) a function of $\sum_f \varepsilon_{1,f} \eta \left( A_{ff} \tilde{m}_{1,f}, \tilde{m}_1 \right)$.

We have then investigated the bounds on the flavour-specific decay asymmetries $\varepsilon_{1,f}$. In the type I seesaw case, it is known that the bound on the flavour-specific asymmetries $\varepsilon_{1,f}$ is substantially relaxed [24] compared to the bound on $\varepsilon_1 = \sum_f \varepsilon_{1,f}$ [44] in the case of a quasi-degenerate spectrum of light neutrinos. For experimentally allowed light neutrino masses below about 0.4 eV, there is no longer a bound on the neutrino mass scale from the requirement of successful thermal leptogenesis. In the type II seesaw case, we have derived the bound on the flavour-specific decay asymmetries $\varepsilon_{1,f} = \varepsilon_{1,f}^{\text{I}} + \varepsilon_{1,f}^{\text{II}}$, which turns out to be identical to the bound on the total decay asymmetry $\varepsilon_1 = \sum_f \varepsilon_{1,f}$. We have compared the bound on the flavour-specific decay asymmetries in type I and type II scenarios, and found that while the type II bound increases with the neutrino mass scale, the type I bound decreases (for experimentally allowed light neutrino masses below about 0.4 eV). The relaxed bound on $\varepsilon_{1,f}$ (figure 5) leads to a lower bound on the mass of the lightest right-handed neutrino $M_{R1}$ in the type II seesaw scenario (figure 7), which decreases when the neutrino mass scale increases. Furthermore, it leads to a relaxed lower bound on the reheat temperature $T_{RH}$ of the early universe (figure 8), which helps to improve consistency of thermal leptogenesis with upper bounds on $T_{RH}$ in some supergravity models. This is in contrast to the type I seesaw scenario, where the lower bound on $T_{RH}$ from thermal leptogenesis increases with increasing neutrino mass scale. Constraints on $T_{RH}$ can therefore imply constraints on the mass scale of the light neutrinos also in flavour-dependent type I leptogenesis, although a general bound is absent.

We have furthermore argued that these relaxed bounds on $\varepsilon_{1,f}$, $M_{R1}$ and $T_{RH}$ in the type II case can be nearly saturated in an elegant way in classes of so-called “type-II-upgraded” seesaw models [8], where the type II contribution to the neutrino mass matrix is proportional to the unit matrix (enforced e.g. by an SO(3) flavour symmetry or by one of its non-Abelian subgroups). One interesting application of these type II seesaw scenarios is that the consistency of thermal leptogenesis with unified theories of flavour is improved compared to the type I seesaw case. This effect, investigated in the flavour-independent approximation in [17], is also present analogously in the flavour-dependent treatment of leptogenesis. The reason is that if the type II contribution ($\propto \mathbb{1}$) dominates, the decay asymmetries $\varepsilon_{1,f}$ become approximately equal and the estimate for the produced baryon asymmetry is similar to the flavour-independent case. Nevertheless,
an accurate analysis of leptogenesis in this scenario requires careful inclusion of the flavour-dependent effects. In many applications and realisations of type II leptogenesis in specific models of fermion masses and mixings (see e.g. [18]), flavour-dependent effects may substantially change the results and they therefore have to be taken into account.

In summary, type II leptogenesis provides a well-motivated generalisation of the conventional scenario of leptogenesis in the type I seesaw framework. We have argued that flavour-dependent effects have to be included in type II leptogenesis, and can change predictions of existing models as well as open up new possibilities for for successful models of leptogenesis. Comparing bounds on $\varepsilon_1$, $M_{R1}$ and $T_{RH}$ in flavour-dependent thermal type I and type II leptogenesis scenarios, we have shown that while type II leptogenesis becomes more efficient for larger mass scale of the light neutrinos, and the bounds become relaxed, leptogenesis within the type I seesaw framework becomes more constrained.

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References

[1] M. Fukugita and T. Yanagida, Phys. Lett. 174B (1986), 45.
[2] D. N. Spergel et al., [arXiv:astro-ph/0603449].
[3] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, Phys. Lett. B 155 (1985) 36.
[4] P. Minkowski, Phys. Lett. B 67 (1977) 421; M. Gell-Mann, P. Ramond and R. Slansky, in Complex Spinors and Unified Theories eds. P. Van. Nieuwenhuizen and D. Z. Freedman, Supergravity (North-Holland, Amsterdam, 1979), p.315 [Print-80-0576 (CERN)]; T. Yanagida, in Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe, eds. O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979), p.95; S. L. Glashow, in Quarks and Leptons, eds. M. Lévy et al. (Plenum Press, New York, 1980), p.687; R. N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 44 (1980) 912.
[5] For a review, see e.g.: W. Buchmuller, R. D. Peccei and T. Yanagida, Ann. Rev. Nucl. Part. Sci. 55 (2005) 311.
[6] A.D. Sakharov. JETP Lett., 5:24, 1967.
[7] R. Barbieri, D. V. Nanopolous, G. Morchio and F. Strocchi, Phys. Lett. B 90 (1980) 91; R. E. Marshak and R. N. Mohapatra, *Invited talk given at Orbis Scientiae, Coral Gables, Fla., Jan. 14-17, 1980*, VPI-HEP-80/02; T. P. Cheng and L. F. Li, Phys. Rev. D 22 (1980) 2860; M. Magg and C. Wetterich, Phys. Lett. B 94 (1980) 61; G. Lazarides, Q. Shafi and C. Wetterich, Nucl. Phys. B181 (1981) 287; J. Schechter and J. W. F. Valle, Phys. Rev. D 22 (1980) 2227; R. N. Mohapatra and G. Senjanović, Phys. Rev. D 23 (1981) 165.

[8] S. Antusch and S. F. King, Nucl. Phys. B 705 (2005) 239.

[9] P. J. O’Donnell and U. Sarkar, Phys. Rev. D 49 (1994), 2118–2121.

[10] E. Ma and U. Sarkar, Phys. Rev. Lett. 80 (1998), 5716–5719.

[11] T. Hambye, E. Ma, and U. Sarkar, Nucl. Phys. B 602 (2001), 23–38.

[12] T. Hambye, M. Raidal and A. Strumia, Phys. Lett. B 632 (2006) 667.

[13] G. Lazarides and Q. Shafi, Phys. Rev. D 58 (1998), 071702.

[14] E. J. Chun and S. K. Kang, Phys. Rev. D 63 (2001) 097902.

[15] T. Hambye and G. Senjanovic, Phys. Lett. B 582, 73 (2004).

[16] S. Antusch and S. F. King, Phys. Lett. B 597 (2004) 199.

[17] S. Antusch and S. F. King, JHEP 0601 (2006) 117.

[18] A. S. Joshipura, E. A. Paschos, and W. Rodejohann, Nucl. Phys. B 611 (2001), 227–238; JHEP 08 (2001), 029; W. Rodejohann, Phys. Lett. B 542 (2002), 100–110; Phys. Rev. D 70 (2004) 073010; P. h. Gu and X. j. Bi, Phys. Rev. D 70 (2004) 063511, hep-ph/0405092; N. Sahu and S. Uma Sankar, Phys. Rev. D 71 (2005) 013006; W. l. Guo, Phys. Rev. D 70 (2004) 053009; N. Sahu and S. Uma Sankar, Nucl. Phys. B 724 (2005) 329; B. Dutta, Y. Mimura and R. N. Mohapatra, Phys. Rev. D 72 (2005) 075009; K. S. Babu, A. Bachri and H. Aissaoui, Nucl. Phys. B 738 (2006) 76; K. Kiers, M. Assis, D. Simons, A. A. Petrov and A. Soni, Phys. Rev. D 73 (2006) 033009; P. Hosteins, S. Lavignac and C. A. Savoy, Nucl. Phys. B 755 (2006) 137; N. Sahu and U. Sarkar, Phys. Rev. D 74 (2006) 093002; P. H. Gu, H. Zhang and S. Zhou, Phys. Rev. D 74 (2006) 076002; F. R. Joaquim and A. Rossi, Nucl. Phys. B 765 (2007) 71; E. J. Chun and S. Scopel, Phys. Rev. D 75 (2007) 023508; M. C. Chen and K. T. Mahanthappa, Phys. Rev. D 75 (2007) 015001; R. N. Mohapatra and H. B. Yu, Phys. Lett. B 644 (2007) 346; A. K. Sarma, H. Zeen Devi and N. N. Singh, Nucl. Phys. B 765 (2007) 142; E. K. Akhmedov, M. Blennow, T. Hallgren, T. Konstandin and T. Ohlsson, arXiv:hep-ph/0612194; Y. Wakabaysi, arXiv:hep-ph/0702261.

[19] R. Barbieri, P. Creminelli, A. Strumia and N. Tetradis, Nucl. Phys. B 575 (2000) 61 [arXiv:hep-ph/9911315].
[20] T. Endoh, T. Morozumi and Z. h. Xiong, Prog. Theor. Phys. 111 (2004) 123 [arXiv:hep-ph/0308276]; T. Fujihara, S. Kaneko, S. Kang, D. Kimura, T. Morozumi and M. Tanimoto, Phys. Rev. D 72 (2005) 016006 [arXiv:hep-ph/0505076].

[21] O. Vives, Phys. Rev. D 73 (2006) 073006 [arXiv:hep-ph/0512160].

[22] A. Abada, S. Davidson, F. X. Josse-Michaux, M. Losada and A. Riotto, JCAP 0604 (2006) 004 [arXiv:hep-ph/0601083];

[23] E. Nardi, Y. Nir, E. Roulet and J. Racker, JHEP 0601 (2006) 164 [arXiv:hep-ph/0601084].

[24] A. Abada, S. Davidson, A. Ibarra, F. X. Josse-Michaux, M. Losada and A. Riotto, [arXiv:hep-ph/0605281]

[25] S. Blanchet and P. Di Bari, [arXiv:hep-ph/0607330].

[26] S. Antusch, S. F. King and A. Riotto, JCAP 0611 (2006) 011 [arXiv:hep-ph/0609038].

[27] S. Pascoli, S. T. Petcov and A. Riotto, [arXiv:hep-ph/0609125]

[28] G. C. Branco, R. Gonzalez Felipe and F. R. Joaquim, Phys. Lett. B 645 (2007) 432 [arXiv:hep-ph/0609297].

[29] S. Antusch and A. M. Teixeira, JCAP 0702 (2007) 024 [arXiv:hep-ph/0611232].

[30] S. Uhlig, [arXiv:hep-ph/0612262]

[31] S. Pascoli, S. T. Petcov and A. Riotto, [arXiv:hep-ph/0611338]

[32] S. Blanchet, P. Di Bari and G. G. Raffelt, [arXiv:hep-ph/0611337].

[33] A. De Simone and A. Riotto, JCAP 0702 (2007) 005 [arXiv:hep-ph/0611357].

[34] G. Engelhard, Y. Grossman, E. Nardi and Y. Nir, [arXiv:hep-ph/0612187]

[35] A. De Simone and A. Riotto, [arXiv:hep-ph/0703175]

[36] F. X. Josse-Michaux and A. Abada, [arXiv:hep-ph/0703084].

[37] T. Shindou and T. Yamashita, [arXiv:hep-ph/0703183].

[38] For recent works where flavour effects are taken into account, see e.g.: A. Pilaftsis and T. E. J. Underwood, Phys. Rev. D 72 (2005) 113001 [arXiv:hep-ph/0506107]; G. C. Branco, A. J. Buras, S. Jager, S. Uhlig and A. Weiler, [arXiv:hep-ph/0609067].

[39] L. Covi, E. Roulet, and F. Vissani, Phys. Lett. B384 (1996), 169–174.

[40] E. Ma, Phys. Rev. Lett. 81 (1998), 1171–1174.
[41] S. Antusch, J. Kersten, M. Lindner, and M. Ratz, Nucl. Phys. B674 (2003), 401–433.

[42] S. Antusch and S. F. King, Phys. Lett. B 591 (2004) 104 [arXiv:hep-ph/0403053].

[43] W. Buchmüller, P. Di Bari, and M. Plümacher, Nucl. Phys. B665 (2003), 445–468.

[44] S. Davidson and A. Ibarra, Phys. Lett. B535 (2002), 25–32.

[45] G. F. Giudice, A. Notari, M. Raidal, A. Riotto, and A. Strumia, Nucl. Phys. B 685, 89 (2004).