Flavor-dependent type II leptogenesis

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We reanalyze leptogenesis via the out-of-equilibrium decay of the lightest right-handed neutrino in type II seesaw scenarios, taking into account flavor-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 flavor-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 flavor-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, flavor-dependent effects cannot be ignored when dealing with leptogenesis in type II seesaw models.

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I. 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 generalized to a type II seesaw [7], where an additional direct mass term m_{LL}^{II} 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 flavor 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 [9–12]. In the first case, there are additional one-loop diagrams where virtual triplets are running in the loop [9,13–16]. In the following, we focus on this possibility, and assume hierarchical right-handed neutrino masses (and that the triplets are heavier than ν_R^1). In this limit, to a good approximation the decay asymmetry depends mainly on the low-energy neutrino mass matrix $m_{LL}^\nu = m_{LL}^I + m_{LL}^{II}$ 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 flavor-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 flavor models with respect to thermal leptogenesis [17]. Finally, since the type II contribution typically does not affect 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 realizations of type II leptogenesis in specific models of fermion masses and mixings, see e.g. [18].

In recent years, the impact of flavor in thermal leptogenesis has merited increasing attention [19–38]. In fact, the one-flavor approximation is only rigorously correct when the interactions mediated by the charged lepton Yukawa couplings are out of equilibrium. Below a given

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temperature (e.g. $\mathcal{O}(10^{12}$ GeV) in the SM and $(1 + \tan^2\beta) \times \mathcal{O}(10^{12}$ GeV) in the MSSM), the tau Yukawa coupling comes into equilibrium (later followed by the couplings of the muon and electron). Flavor 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 flavor. In the full computation, the asymmetries in each distinguishable flavor are differently washed out, and appear with distinct weights in the final baryon asymmetry.

Flavor-dependent leptogenesis in the type I seesaw scenario has recently been addressed in detail by several authors. In particular, flavor-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 flavor space have been included in [22,32,33,35] and the treatment has been generalized 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 [29,36] have also been considered including flavor-dependent effects. Strong connections between the low-energy *CP* phases of the U_{MNS} matrix and *CP* violation for flavor-dependent leptogenesis have been shown to emerge in certain classes of neutrino mass models [26] 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 flavor-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 flavor-dependent effects are included [24].

In view of the importance of flavor-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 reanalyze leptogenesis via the out-of-equilibrium decay of the lightest right-handed neutrino in type II seesaw scenarios, taking into account flavor-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.

II. 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 the MSSM, respectively,

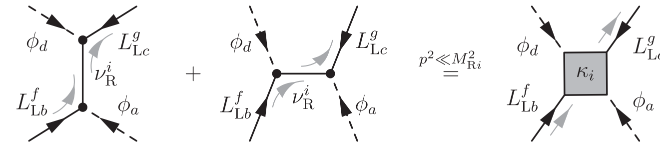

FIG. 1. Generation of the dimension 5 neutrino mass operator in the type I seesaw mechanism.

$$\mathcal{L}_\kappa^{\mathrm{SM}} = \tfrac{1}{4}\kappa_{gf}(\overline{L^{Cg}} \cdot \phi)(L^f \cdot \phi) + \mathrm{H.c.}, \tag{1a}$$

$$\mathcal{L}_\kappa^{\mathrm{MSSM}} = -\tfrac{1}{4}\kappa_{gf}(\hat{L}^g \cdot \hat{H}_{\mathrm{u}})(\hat{L}^f \cdot \hat{H}_{\mathrm{u}})|_{\theta\theta} + \mathrm{H.c.}, \tag{1b}$$

is generated from integrating out the heavy fields. This is illustrated in Figs. 1 and 2. In Eq. (1), the dots indicate the $SU(2)_L$-invariant product, $(\hat{L}^f \cdot \hat{H}_{\mathrm{u}}) = \hat{L}^f_a (i\tau_2)^{ab}(\hat{H}_{\mathrm{u}})_b$, with τ_A $(A \in \{1, 2, 3\})$ being the Pauli matrices. Superfields are marked by hats. After electroweak symmetry breaking, the operators of Eq. (1) lead to Majorana mass terms for the light neutrinos,

$$\mathcal{L}_\nu = -\frac{1}{2} m^\nu_{\mathrm{LL}} \bar{\nu}_{\mathrm{L}} \nu^{\mathrm{C}f}_{\mathrm{L}}, \quad \text{with} \quad m^\nu_{\mathrm{LL}} = -\frac{v_{\mathrm{u}}^2}{2}(\kappa)^*. \tag{2}$$

In the type I seesaw mechanism, it is assumed that only the singlet (right-handed) neutrinos $\nu_{\mathrm{R}i}$ contribute to the neutrino masses. With Y_ν being the neutrino Yukawa matrix in left-right convention,[1] M_{RR} the mass matrix of the right-handed neutrinos, and $v_{\mathrm{u}} = \langle \phi^0 \rangle (= \langle H^0_{\mathrm{u}} \rangle)$ 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

$$m^{\mathrm{I}}_{\mathrm{LL}} = -v_{\mathrm{u}}^2 Y_\nu M^{-1}_{\mathrm{RR}} Y^T_\nu. \tag{3}$$

In the type II seesaw mechanism, the contributions to the neutrino mass matrix from both, right-handed neutrinos $\nu_{\mathrm{R}i}$ and Higgs triplet(s) Δ_{L}, are considered. The additional contribution to the neutrino masses from Δ_{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. Fig. 2). The neutrino mass matrix in the type II seesaw mechanism has the form

$$m^\nu_{\mathrm{LL}} = m^{\mathrm{II}}_{\mathrm{LL}} + m^{\mathrm{I}}_{\mathrm{LL}} = m^{\mathrm{II}}_{\mathrm{LL}} - v_{\mathrm{u}}^2 Y_\nu M^{-1}_{\mathrm{RR}} Y^T_\nu, \tag{4}$$

[1] The neutrino Yukawa matrix corresponds to $-(Y_\nu)_{fi} \times (L^f \cdot \phi)\nu^i_{\mathrm{R}}$ in the Lagrangian of the SM and, analogously, to $(Y_\nu)_{fi}(\hat{L}^f \cdot \hat{H}_{\mathrm{u}})\hat{\nu}^{\mathrm{C}i}$ in the superpotential of the MSSM (see [16] for further details).

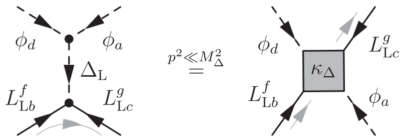

FIG. 2. Extra diagram generating the dimension 5 neutrino mass operator in the type II seesaw mechanism from a $SU(2)_L$-triplet Higgs field.

where m^{II}_{LL} is the additional term from the Higgs triplet(s). In left-right symmetric unified theories, the generic size of both seesaw contributions m^{I}_{LL} and m^{II}_{LL} is $\mathcal{O}(v_u^2/v_{B-L})$ 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)).

III. BARYOGENESIS VIA FLAVOR-DEPENDENT LEPTOGENESIS

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

The relevance of the flavor-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 $(1+\tan^2\beta)\times 10^5$ GeV and $(1+\tan^2\beta)\times 10^9$ GeV, the μ and τ Yukawa couplings are in thermal equilibrium and all flavors in the Boltzmann equations are to be treated separately. For $\tan\beta = 30$, this applies for temperatures below about 10^{12} GeV and above 10^8 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 GeV and 10^{12} GeV, only the τ Yukawa coupling is in equilibrium and is treated separately in the Boltzmann equations, whereas μ and e flavors are indistinguishable. A discussion of the temperature regimes in the SM and MSSM, where flavor is important, can be found, e.g., in [26].

We now briefly review the estimation of the produced baryon asymmetry in flavor-dependent leptogenesis.[2] For definiteness, we focus on the temperature range where all flavors are to be treated separately. In the following discussion of thermal type II leptogenesis, we will assume that the mass M_{Δ_L} of the triplet(s) is much larger than M_{R1}. In this limit, the flavor-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 ν^1_R and $\tilde{\nu}^1_R$ give rise to flavor-dependent asymmetries in the (s)lepton sector, which are then partly transformed via sphaleron conversion into a baryon asymmetry Y_B.[3] The final baryon asymmetry can be calculated as

$$Y_B^{SM} = \frac{12}{37}\sum_f Y^{SM}_{\Delta_f}, \tag{5}$$

$$Y_B^{MSSM} = \frac{10}{31}\sum_f \hat{Y}^{MSSM}_{\Delta_f}, \tag{6}$$

where $\hat{Y}_{\Delta_f} \equiv Y_B/3 - Y_{L_f}$ are the total (particle and sparticle) $B/3 - L_f$ asymmetries, with Y_{L_f} the lepton number densities in the flavor $f = e, \mu, \tau$. The asymmetries $\hat{Y}^{MSSM}_{\Delta_f}$ and $Y^{SM}_{\Delta_f}$, 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 parametrize the produced asymmetries in terms of flavor-specific efficiency factors η_f and decay asymmetries $\varepsilon_{1,f}$ as

$$Y^{SM}_{\Delta_f} = \eta^{SM}_f \varepsilon_{1,f} Y^{eq}_{\nu^1_R}, \tag{7}$$

$$\hat{Y}^{MSSM}_{\Delta_f} = \eta^{MSSM}_f [\tfrac{1}{2}(\varepsilon_{1,f} + \varepsilon_{1,\tilde{f}}) Y^{eq}_{\nu^1_R} + \tfrac{1}{2}(\varepsilon_{\tilde{1},f} + \varepsilon_{\tilde{1},\tilde{f}}) Y^{eq}_{\tilde{\nu}^1_R}]. \tag{8}$$

$Y^{eq}_{\nu^1_R}$ and $Y^{eq}_{\tilde{\nu}^1_R}$ are the number densities of the neutrino and sneutrino for $T \gg M_1$ if they were in thermal equilibrium, normalized with respect to the entropy density. In the Boltzmann approximation, they are given by $Y^{eq}_{\nu^1_R} \approx Y^{eq}_{\tilde{\nu}^1_R} \approx 45/(\pi^4 g_*)$. g^* is the effective number of degrees of freedom, which amounts 106.75 in the SM and 228.75 in the MSSM.

$\varepsilon_{1,f}$, $\varepsilon_{1,\tilde{f}}$, $\varepsilon_{\tilde{1},f}$, and $\varepsilon_{\tilde{1},\tilde{f}}$ 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

$$\varepsilon_{1,f} = \frac{\Gamma_{\nu^1_R L_f} - \Gamma_{\nu^1_R \bar{L}_f}}{\sum_f(\Gamma_{\nu^1_R L_f} + \Gamma_{\nu^1_R \bar{L}_f})}, \qquad \varepsilon_{1,\tilde{f}} = \frac{\Gamma_{\nu^1_R \tilde{L}_f} - \Gamma_{\nu^1_R \tilde{L}^*_f}}{\sum_f(\Gamma_{\nu^1_R \tilde{L}_f} + \Gamma_{\nu^1_R \tilde{L}^*_f})},$$
$$\varepsilon_{\tilde{1},f} = \frac{\Gamma_{\tilde{\nu}^{*1}_R L_f} - \Gamma_{\tilde{\nu}^1_R \bar{L}_f}}{\sum_f(\Gamma_{\tilde{\nu}^{*1}_R L_f} + \Gamma_{\tilde{\nu}^1_R \bar{L}_f})}, \qquad \varepsilon_{\tilde{1},\tilde{f}} = \frac{\Gamma_{\tilde{\nu}^1_R \tilde{L}_f} - \Gamma_{\tilde{\nu}^{*1}_R \tilde{L}^*_f}}{\sum_f(\Gamma_{\tilde{\nu}^1_R \tilde{L}_f} + \Gamma_{\tilde{\nu}^{*1}_R \tilde{L}^*_f})}. \tag{9}$$

[2]For a discussion of approximations which typically enter these estimates, and which also apply to our discussion, see e.g. Sec. 3.1.3 in [29].

[3]In the following, Y will always be used for quantities which are normalized to the entropy density s. The quantities normalized with respect to the photon ratio can be obtained using the relation $s/n_\gamma \approx 7.04k$.

The flavor-dependent efficiency factors η_f in the SM and in the MSSM are defined by Eqs. (7) and (8), respectively. As stated above, we assume that the mass M_{Δ_L} of the triplet(s) is much larger than M_{R1}. In this limit, the efficiencies for flavor-dependent thermal leptogenesis in the type I and type II frameworks are mainly determined by the properties of ν_R^1, which means, in particular, that the flavor-dependent efficiencies 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 ν_{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 flavor-dependent Boltzmann equations, which can be found for the SM in [19,22–24] and for the MSSM in [26,29]. In general, the flavor-dependent efficiencies depend strongly on the washout parameters $\tilde{m}_{1,f}$ for each flavor, and on the total washout parameter $\tilde{m}_1$, which are defined as

$$\tilde{m}_{1,f} = \frac{v_u^2 (Y_\nu^\dagger Y_\nu)_{11}}{M_{R1}}, \qquad \tilde{m}_1 = \sum_f \tilde{m}_{1,f}. \tag{10}$$

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

$$K_f = \frac{\tilde{m}_{1,f}}{m^*}, \qquad K = \sum_f K_f, \tag{11}$$

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

The most relevant difference between the flavor-independent approximation and the correct flavor-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 flavor-independent approximation where the lepton flavor is neglected in the Boltzmann equations. In other words, in the flavor-independent approximation the total baryon asymmetry is a function of $(\sum_f \varepsilon_{1,f}) \times \eta^{\mathrm{ind}}(\sum_g K_g)$. In the correct flavor treatment the baryon asymmetry is (approximately) a function of $\sum_f \varepsilon_{1,f}\eta(A_{ff}K_f, K)$. From this, it is already clear that flavor-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 η_f (which depend mainly on $\tilde{m}_{1,f}$ and $\tilde{m}_1$ (or K_f and K)). While the efficiency factors

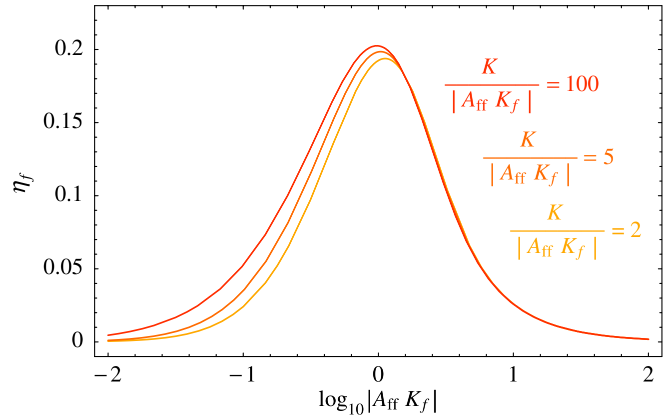

FIG. 3 (color online). Flavor-dependent efficiency factor $\eta(A_{ff}K_f, K)$ in the MSSM as a function of $A_{ff}K_f$, for fixed values of $K/|A_{ff}K_f| = 2$, 5, and 100, obtained from solving the flavor-dependent Boltzmann equations in the MSSM with zero initial abundance of right-handed (s)neutrinos (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 $|A_{ff}|$ of $\mathcal{O}(1)$. 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 flavor-dependent efficiency factors for different $K/|A_{ff}K_f|$ is that the total baryon asymmetry is the sum of each individual lepton asymmetries, which is weighted by the corresponding 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.

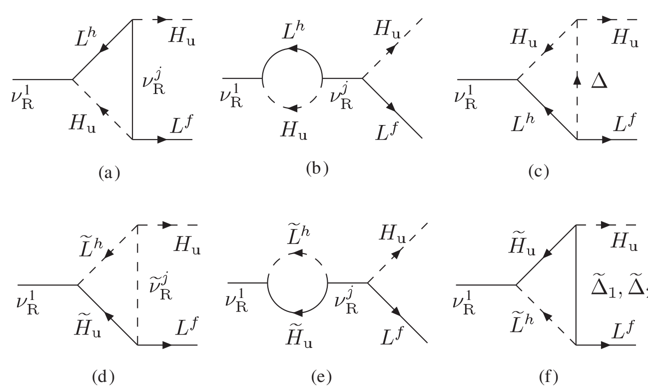

FIG. 4. Loop diagrams in the MSSM which contribute to the decay $\nu_R^1 \to L_a^f H_{ub}$ for the case of a type II seesaw mechanism where the direct mass term for the neutrinos stems from the induced vev of a Higgs triplet. In diagram (f), $\tilde{\Delta}_1$ and $\tilde{\Delta}_2$ are the mass eigenstates corresponding to the superpartners of the $SU(2)_L$-triplet scalar fields Δ 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 ϕ.

IV. DECAY ASYMMETRIES

A. 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 $\nu_R^1 \to L_a^f H_{ub}$ are shown in Fig. 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 flavor yields

$$\varepsilon_{1,f}^{(a)} = \frac{1}{8\pi} \frac{\sum_{j\neq 1} \mathrm{Im}[(Y^\dagger)_{1f}(Y_\nu^\dagger Y_\nu)_{1j}(Y^T)_{jf}]}{(Y_\nu^\dagger Y_\nu)_{11}} \sqrt{x_j}\left[1-(1+x_j)\ln\left(\frac{x_j+1}{x_j}\right)\right], \tag{12a}$$

$$\varepsilon_{1,f}^{(b)} = \frac{1}{8\pi} \frac{\sum_{j\neq 1} \mathrm{Im}[(Y^\dagger)_{1f}(Y_\nu^\dagger Y_\nu)_{1j}(Y^T)_{jf}]}{(Y_\nu^\dagger Y_\nu)_{11}} \sqrt{x_j}\left[\frac{1}{1-x_j}\right], \tag{12b}$$

$$\varepsilon_{1,f}^{(c)} = -\frac{3}{8\pi}\frac{M_{R1}}{v_u^2} \frac{\sum_g \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{II})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}} y\left[-1+y\ln\left(\frac{y+1}{y}\right)\right], \tag{12c}$$

$$\varepsilon_{1,f}^{(d)} = \frac{1}{8\pi} \frac{\sum_{j\neq 1} \mathrm{Im}[(Y^\dagger)_{1f}(Y_\nu^\dagger Y_\nu)_{1j}(Y^T)_{jf}]}{(Y_\nu^\dagger Y_\nu)_{11}} \sqrt{x_j}\left[-1+x_j\ln\left(\frac{x_j+1}{x_j}\right)\right], \tag{12d}$$

$$\varepsilon_{1,f}^{(e)} = \frac{1}{8\pi} \frac{\sum_{j\neq 1} \mathrm{Im}[(Y^\dagger)_{1f}(Y_\nu^\dagger Y_\nu)_{1j}(Y^T)_{jf}]}{(Y_\nu^\dagger Y_\nu)_{11}} \sqrt{x_j}\left[\frac{1}{1-x_j}\right], \tag{12e}$$

$$\varepsilon_{1,f}^{(f)} = -\frac{3}{8\pi}\frac{M_{R1}}{v_u^2} \frac{\sum_g \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{II})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}} y\left[1-(1+y)\ln\left(\frac{y+1}{y}\right)\right], \tag{12f}$$

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 Eq. (12) they have been generalized to the flavor-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

$$\varepsilon_{1,f}^{\mathrm{SM,II}} = \varepsilon_{1,f}^{(c)}, \tag{13a}$$

$$\varepsilon_{1,f}^{\mathrm{MSSM,II}} = \varepsilon_{1,f}^{(c)} + \varepsilon_{1,f}^{(f)}. \tag{13b}$$

In the MSSM, we furthermore obtain

$$\varepsilon_{1,f}^{\mathrm{MSSM,II}} = \varepsilon_{1,\tilde{f}}^{\mathrm{MSSM,II}} = \varepsilon_{\tilde{1},f}^{\mathrm{MSSM,II}} = \varepsilon_{\tilde{1},\tilde{f}}^{\mathrm{MSSM,II}}. \tag{14}$$

The results corresponding to the diagrams (a), (b), (d), and (e) which contribute to $\varepsilon_1^{\mathrm{I}}$ 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 MSSM are

$$\varepsilon_{1,f}^{\mathrm{SM,I}} = \varepsilon_{1,f}^{(a)} + \varepsilon_{1,f}^{(b)}, \tag{15a}$$

$$\varepsilon_{1,f}^{\mathrm{MSSM,I}} = \varepsilon_{1,f}^{(a)} + \varepsilon_{1,f}^{(b)} + \varepsilon_{1,f}^{(d)} + \varepsilon_{1,f}^{(e)}. \tag{15b}$$

Again, in the MSSM, the remaining decay asymmetries are equal to $\varepsilon_{1,f}^{\mathrm{MSSM,I}}$:

$$\varepsilon_{1,f}^{\mathrm{MSSM,I}} = \varepsilon_{1,\tilde{f}}^{\mathrm{MSSM,I}} = \varepsilon_{\tilde{1},f}^{\mathrm{MSSM,I}} = \varepsilon_{\tilde{1},\tilde{f}}^{\mathrm{MSSM,I}}. \tag{16}$$

Finally, the total decay asymmetries from the decay of ν_R^1 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

$$\varepsilon_{1,f}^{\mathrm{SM}} = \varepsilon_{1,f}^{\mathrm{SM,I}} + \varepsilon_{1,f}^{\mathrm{SM,II}}, \tag{17}$$

$$\varepsilon_{1,f}^{\mathrm{MSSM}} = \varepsilon_{1,f}^{\mathrm{MSSM,I}} + \varepsilon_{1,f}^{\mathrm{MSSM,II}}. \tag{18}$$

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

$$\frac{\sum_{j\neq 1} \mathrm{Im}[(Y^\dagger)_{1f}(Y_\nu^\dagger Y_\nu)_{1j}(Y^T)_{jf}]}{8\pi(Y_\nu^\dagger Y_\nu)_{11}}\frac{1}{\sqrt{x_j}} = -\frac{M_{R1}}{v_u^2}\frac{\sum_g \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{\mathrm{I}})_{fg}]}{8\pi(Y_\nu^\dagger Y_\nu)_{11}}. \tag{19}$$

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_Δ, we obtain the simple results for the flavor-specific decay asymmetries $\varepsilon_{1,f}^{SM}$ and $\varepsilon_{1,f}^{MSSM}$ [16]

$$\varepsilon_{1,f}^{SM} = \frac{3}{16\pi} \frac{M_{R1}}{v_u^2} \frac{\sum_{fg} \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{I} + m_{LL}^{II})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}}, \tag{20a}$$

$$\varepsilon_{1,f}^{MSSM} = \frac{3}{8\pi} \frac{M_{R1}}{v_u^2} \frac{\sum_{fg} \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{I} + m_{LL}^{II})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}}. \tag{20b}$$

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.

B. 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 not produced nonthermally 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 [16].

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

$$(m_{LL}^\nu)_{fg} = -\frac{v_u^2}{2}[2(Y_\nu)_{f1}M_{R1}^{-1}(Y_\nu^T)_{1g} + (\kappa'^*)_{fg}]. \tag{21}$$

κ' includes type I contributions from the heavier right-handed neutrinos, plus any additional (type II) contributions from heavier particles. Examples for realizations of the neutrino mass operator can be found, e.g., in [40].

At M_{R1}, the minimal effective field theory extension of the SM (MSSM) for leptogenesis includes the effective neutrino mass operator κ' plus one right-handed neutrino ν_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 $(Y_\nu)_{f1} \times (\hat{L}^f \cdot \hat{H}_u)\hat{\nu}^{C1}$ 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 ν_{R1} (and $\tilde{\nu}_{R1}$) with the loop diagrams containing the effective operator, shown in Fig. 5. In the SM, we obtain the simple result [16] for the flavor-specific effective decay asymmetries (corresponding to diagram (a) of Fig. 5)

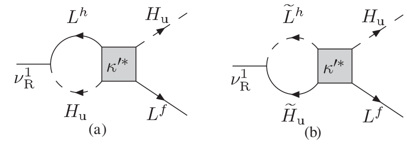

FIG. 5. Loop diagrams contributing to the decay asymmetry via the decay $\nu_R^1 \to L_a^f H_{ub}$ in the MSSM with a (lightest) right-handed neutrino ν_R^1 and a neutrino mass matrix determined by κ' [16]. Further contributions to the generated baryon asymmetry stem from the decay of ν_R^1 into slepton and Higgsino and from the decays of the sneutrino $\tilde{\nu}_R^1$. With H_u renamed to the SM Higgs, the first diagram contributes in the extended SM.

$$\varepsilon_{1,f}^{SM} = \frac{3}{16\pi} \frac{M_{R1}}{v_u^2} \frac{\sum_{g} \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{\nu})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}}. \tag{22}$$

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

$$\varepsilon_{1,f}^{MSSM} = \frac{3}{8\pi} \frac{M_{R1}}{v_u^2} \frac{\sum_{g} \mathrm{Im}[(Y_\nu^*)_{f1}(Y_\nu^*)_{g1}(m_{LL}^{\nu})_{fg}]}{(Y_\nu^\dagger Y_\nu)_{11}}. \tag{23}$$

Explicit calculation furthermore yields

$$\varepsilon_{1,f}^{MSSM} = \varepsilon_{1,\tilde{f}}^{MSSM} = \varepsilon_{\tilde{1},f}^{MSSM} = \varepsilon_{\tilde{1},\tilde{f}}^{MSSM}. \tag{24}$$

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

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 Eq. (12) [16], in the limit $M_\Delta \gg M_{R1}$.

V. 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 $|\varepsilon_1^{SM}| = |\sum_f \varepsilon_{1,f}^{SM}|$ and $|\varepsilon_1^{MSSM}| = |\sum_f \varepsilon_{1,f}^{MSSM}|$, have been derived in [16]. For the flavor-specific decay asymmetries $\varepsilon_{1,f}^{SM}$ and $\varepsilon_{1,f}^{MSSM}$, the bounds can readily be obtained as

$$|\varepsilon_{1,f}^{SM}| \le \frac{3}{16\pi}\frac{M_{R1}}{v_u^2} m_{max}^\nu, \qquad |\varepsilon_{1,f}^{MSSM}| \le \frac{3}{8\pi}\frac{M_{R1}}{v_u^2} m_{max}^\nu. \tag{25}$$

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. Fig. 4 of [41]).

A situation where an almost maximal baryon asymmetry is generated by thermal leptogenesis can be realized, 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 flavors approximately take its optimal value. Classes of type II seesaw models, where this can be accommodated, have been considered in [8,17,42]. 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) flavor symmetry or by one of its non-Abelian subgroups). From Eq. (20), one can readily see that if the type II contribution ($\propto \mathbb{1}$) dominates the neutrino mass matrix m_{LL}^{ν}, and if $(Y_\nu)_{f1}$ are approximately equal for all flavors $f = 1, 2, 3$ and chosen such that the resulting $\tilde{m}_{1,f}$ are approximately equal to m^*, we have realized $\eta_f \approx \eta_{max}$ for all flavors and simultaneously nearly saturated the bound for the total decay asymmetry.[4]

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

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

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 flavor-specific decay asymmetries are constrained by [24]

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

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 flavor-independent washout parameter has the lower bound [43]

$$\tilde{m}_1 \geq m_{min}^{\nu}, \quad (28)$$

with $m_{min}^{\nu} = \min(m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$. On the contrary, in the type I and type II seesaw, the flavor-dependent washout parameters $\tilde{m}_{1,f}$ are generically not constrained. Note that in the flavor-independent approximation, Eq. (28) leads to a dramatically more restrictive bound on $\varepsilon_1 = \sum_f \varepsilon_{1,f}$ [44] for quasidegenerate light neutrino masses, and finally even to a bound on the neutrino mass scale [43]. This can be understood from the fact that for $\tilde{m}_1 \gg m^*$ in the flavor-independent approximation, washout effects strongly reduce the efficiency of thermal leptogenesis. Similarly, in the flavor-dependent treatment, $\tilde{m}_{1,f} \gg m^*$ would lead to a strongly reduced efficiency for this specific flavor. This strong washout for quasidegenerate light neutrinos can be avoided in flavor-dependent type I leptogenesis, and $\tilde{m}_{1,f} \approx m^*$ can realize a nearly optimal scenario regarding washout (c.f. Fig. 3). However, we see from Eq. (27) that the decay asymmetries in this case are reduced by a factor of $(m^*/m_{min}^{\nu})^{1/2}$ when compared to the optimal value, leading to a reduced baryon asymmetry. On the other hand, realizing nearly optimal $\varepsilon_{1,f}$ requires $\tilde{m}_{1,f} \approx \tilde{m}_1 \geq m_{min}^{\nu}$, leading to large washout effects for quasidegenerate light neutrinos and even to a more strongly suppressed generation of baryon asymmetry (c.f. Fig. 3). As a consequence, increasing the neutrino mass scale increases the

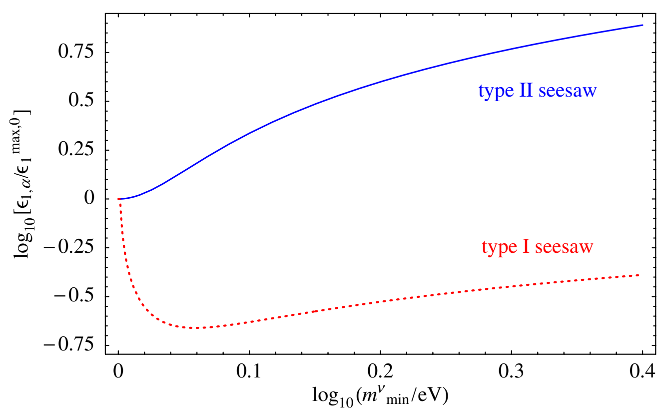

FIG. 6 (color online). Bound on the decay asymmetry $\varepsilon_{1,f}$ in type II leptogenesis (solid blue line) and type I leptogenesis (dotted red line) as a function of the mass of the lightest neutrino $m_{min}^{\nu} := \min(m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$ in type I and type II seesaw scenarios (see also [29]). The washout parameter $|A_{ff}|\tilde{m}_{1,f}$ is fixed to m^* (close to optimal), and the asymmetry is normalized to $\varepsilon_1^{max,0} = 3M_{R1}(\Delta m_{31}^2)_{1/2}/(16\pi v_u^2)$, 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.

[4]We further note that the bound for one of the flavor-specific decay asymmetries can be nearly saturated in this scenario if, for instance, $(Y_\nu)_{21} \approx (Y_\nu)_{31} \approx 0$.

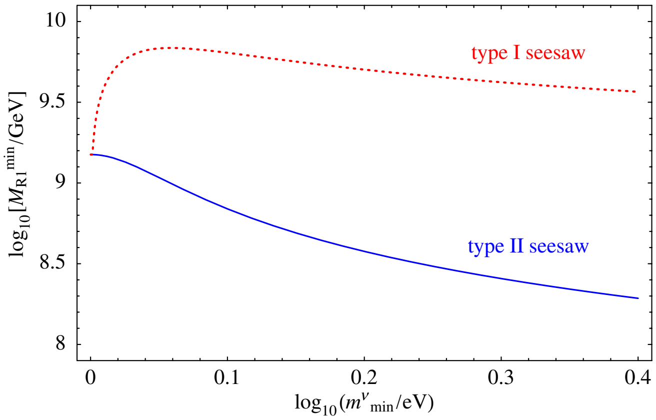

FIG. 7 (color online). Lower bound on M_{R1} in type II leptogenesis (solid blue line) and type I leptogenesis (dotted red line) as a function of the mass of the lightest neutrino $m^{\nu}_{min} := \min(m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$. For definiteness, the MSSM with $\tan\beta = 30$ has been considered as an example.

lower bound on M_{R1} (also in the presence of flavor-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 quasidegenerate light neutrino masses either by a factor $(m^*/m^{\nu}_{min})^{1/2}$ in the decay asymmetries or by a nonoptimal washout parameter much larger than m^* (or $K_f \gg 1$, c.f. Fig. 3). The bounds on the decay asymmetries in type I and type II leptogenesis are compared in Fig. 6, where $\tilde{m}_{1,f}$ has been fixed to m^*, close to its optimal value. From Fig. 6 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 Fig. 7. In addition, for the same reason, increasing the neutrino mass scale also relaxes the lower bound on the reheat temperature T_{RH} from the requirement of successful type II leptogenesis. Including reheating in the flavor-dependent Boltzmann equations as in Ref. [29] (for the flavor-independent case, see [45]), we obtain the m^{ν}_{min}-dependent lower bounds on T_{RH} in type I and type II scenarios shown in Fig. 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 $m^{\nu}_{min} = \min(m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$. For instance, with an upper bound $T_{RH} \leq 5 \times 10^9$ GeV, values of m^{ν}_{min} in the approximate range [0.01 eV, 0.32 eV] would be incompatible with leptogenesis in the type I seesaw framework (c.f. Fig. 8).

VI. SUMMARY, DISCUSSION, AND CONCLUSIONS

We have analyzed flavor-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 flavor-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 flavor-dependent regime the decays of the right-handed (s)neutrinos generate asymmetries in each distinguishable flavor (proportional to the flavor-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 η_f) in the final baryon asymmetry.

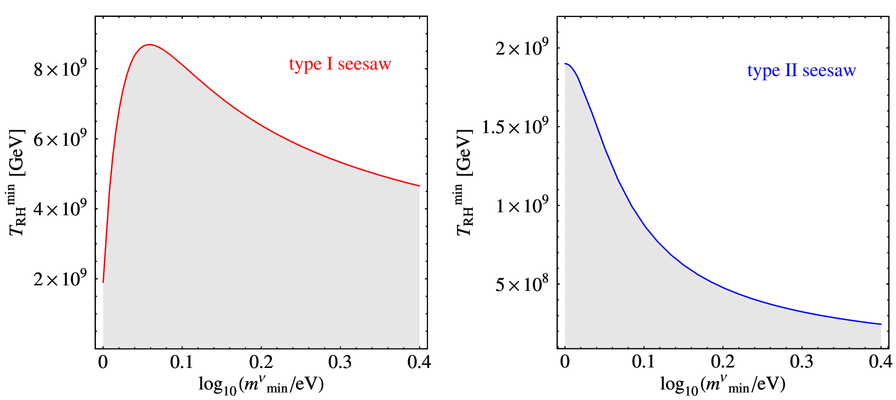

FIG. 8 (color online). Lower bound on the reheat temperature T_{RH} in type I leptogenesis (left panel) and in type II leptogenesis (right panel) as a function of the mass of the lightest neutrino $m^{\nu}_{min} = \min(m_{\nu_1}, m_{\nu_3}, m_{\nu_3})$, in the MSSM with $\tan\beta = 30$. In the gray regions, values of T_{RH} are incompatible with thermal leptogenesis for the corresponding m^{ν}_{min}.

The most important quantities for computing the produced baryon asymmetry are the decay asymmetries $\varepsilon_{1,f}$ and the efficiency factors η_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 flavor-specific efficiency factors η_f, in the limit that the mass M_{Δ_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 generalized the results of [16] to the flavor-dependent case. The most important effects of flavor in leptogenesis are a consequence of the fact that in the flavor-independent approximation the total baryon asymmetry is a function of $(\sum_f \varepsilon_{1,f}) \times \eta^{\text{ind}}(\sum_g \tilde{m}_{1,g})$, whereas in the correct flavor-dependent treatment the baryon asymmetry is (approximately) a function of $\sum_f \varepsilon_{1,f}\eta(A_{ff}\tilde{m}_{1,f}, \tilde{m}_1)$.

We have then investigated the bounds on the flavor-specific decay asymmetries $\varepsilon_{1,f}$. In the type I seesaw case, it is known that the bound on the flavor-specific asymmetries $\varepsilon^{\text{I}}_{1,f}$ is substantially relaxed [24] compared to the bound on $\varepsilon^{\text{I}}_1 = \sum_f \varepsilon^{\text{I}}_{1,f}$ [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 flavor-specific decay asymmetries $\varepsilon_{1,f} = \varepsilon^{\text{I}}_{1,f} + \varepsilon^{\text{II}}_{1,f}$, 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 flavor-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 $\varepsilon_{1,f}$ (Fig. 6) leads to a lower bound on the mass of the lightest right-handed neutrino M_{R1} in the type II seesaw scenario (Fig. 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 (Fig. 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 flavor-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) flavor 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 flavor is improved compared to the type I seesaw case. This effect, investigated in the flavor-independent approximation in [17], is also present analogously in the flavor-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 flavor-independent case. Nevertheless, an accurate analysis of leptogenesis in this scenario requires careful inclusion of the flavor-dependent effects. In many applications and realizations of type II leptogenesis in specific models of fermion masses and mixings (see e.g. [18]), flavor-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 generalization of the conventional scenario of leptogenesis in the type I seesaw framework. We have argued that flavor-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 successful models of leptogenesis. Comparing bounds on $\varepsilon_{1,f}$ M_{R1} and T_{RH} in flavor-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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