

Implications of Flavor Dynamics for Fermion Triplet Leptogenesis

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Abstract

We analyze the importance of flavor effects in models in which leptogenesis proceeds via the decay of Majorana electroweak triplets. We find that depending on the relative strengths of gauge and Yukawa reactions the $B - L$ asymmetry can be sizably enhanced, exceeding in some cases an order of magnitude level. We also discuss the impact that such effects can have for TeV-scale triplets showing that as long as the $B - L$ asymmetry is produced by the dynamics of the lightest such triplet they are negligible, but open the possibility for scenarios in which the asymmetry is generated above the TeV scale by heavier states, possibly surviving the TeV triplet related washouts. We investigate these cases and show how they can be disentangled at the LHC by using Majorana triplet collider observables and, in the case of minimal type III see-saw models even through lepton flavor violation observables.

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1 Introduction

Measurements from light element abundances and the cosmic microwave background radiation allow to determine the cosmic baryon asymmetry, $Y_{\Delta_B} = (8.75 \pm 0.23) \times 10^{-11}$ [1]. The conditions under which this asymmetry can be dynamically generated (baryogenesis) are well known [2] and depending on how they are realized different mechanisms for baryogenesis can be defined [3].

Leptogenesis is a mechanism in which an initial lepton asymmetry is partially reprocessed into a baryon asymmetry by nonperturbative sphaleron interactions [4]. Qualitatively a net non-zero lepton asymmetry can be generated in any framework containing interactions that: (i) break lepton number; (ii) violate CP; (iii) depart from thermal equilibrium at some stage in the cosmic evolution. In principle, these conditions are satisfied in any neutrino mass model in which light neutrinos acquire Majorana masses and therefore in leptogenesis two unrelated puzzles, the origin of the Cosmic baryon asymmetry and of neutrino masses, are linked together. In the standard approach, in which light neutrinos acquire Majorana masses via the type I see-saw model [5], leptogenesis takes place via the CP violating out-of-equilibrium decays of heavy standard model fermionic singlets [6]. This scenario for leptogenesis has been widely discussed in the literature [7] and indeed a lot of progress in the understanding of the generation of the lepton asymmetry within this scenario has been achieved [8, 9, 10, 11, 12, 13, 14, 15].

Given that qualitatively the conditions for producing a lepton asymmetry are also fulfilled in other see-saw realizations (type II [16] and type III [17]) one is tempted to extend the standard type I analysis to these cases, and in fact studies of such scenarios have been considered [18, 19, 20, 21, 22]. In these models the lepton asymmetry is generated via the dynamics of either a scalar (type II) or a Majorana fermion (type III) $SU(2)$ triplet and thus the major difference between these cases and the standard one arises from the fact that both, the scalar and fermion, couple to standard model (SM) gauge bosons. At high temperatures gauge reactions are much faster than the expansion rate of the Universe implying that no asymmetry can be generated at this stage. As the temperature drops, thermalization of the triplet distribution becomes less efficient and depending on the strength of the Yukawa interactions the generation of the lepton asymmetry can proceed either after the decoupling of gauge reactions or after the Yukawa interactions freeze-out [21].

In this paper we focus on leptogenesis in type III see-saw (fermionic triplet leptogenesis). In particular, we analyze the dynamics of lepton flavor in the generation of the lepton asymmetry and the implications that such effects could have for TeV-scale triplets. As in the standard case when the lightest triplet mass (M_{T_1}) is below 10^{12} GeV the charged lepton Yukawa interactions that are sufficiently fast project the lepton and anti-lepton quantum states produced in T_1 decays into their flavor components before they can re-scatter [8, 10, 11]. In the standard case the fermionic singlet dynamics is completely determined by Yukawa reactions and accordingly the impact of flavor effects becomes well pronounced. In contrast, in the fermionic triplet leptogenesis case, since the dynamics of the triplet is not completely determined by its Yukawa interactions, it is not entirely clear whether light flavor effects could yield a sizable enhancement of the final

lepton asymmetry regardless of the strength of the Yukawa reactions. Here we will show that the inclusion of flavor may in general produce an enhancement of the asymmetry, but that relevant effects are possible only when the strength of the triplet Yukawa interactions are such that they are still active when gauge reactions decouple.

In the case of a quasi-degenerate fermionic singlet mass spectrum the resonant enhancement of the CP violating asymmetry allows for successful leptogenesis even when the fermionic singlet masses are well below 1 TeV [23, 24]. In fermionic triplet leptogenesis on the other hand, the efficiency due to gauge reactions strongly depends on M_{T_1} and it is drastically diminished when $\mathcal{O}(M_{T_1}) \sim \text{TeV}$. In the one flavor approximation in ref. [22] it was pointed out that this constraint in conjunction with sphaleron decoupling implies that the correct amount of baryon asymmetry can only be generated for $M_{T_1} \gtrsim 1.6 \text{ TeV}$. As we will show, as long as leptogenesis proceeds through T_1 dynamics, this bound remains valid even when flavor effects are accounted for. However, this does not necessarily imply that the observation of fermionic triplets at the LHC would exclude the possibility of successful leptogenesis. Once flavor effects are taken into account the possibility of high scale leptogenesis remain plausible. As will be discussed, in that case, the $B - L$ asymmetry produced above the TeV-scale (e.g. in T_2 decays) can survive T_1 related washouts only for particular flavor structures. These structures constrain the dynamics of T_1 leading to experimental signatures which could be used to constrain leptogenesis models.

The rest of this paper is organized as follows. In section 2 we discuss the generalities of type III see-saw and review the generation of the $B - L$ asymmetry in the one flavor approximation, we also determine the regions where Yukawa reactions remain active after gauge interactions decouple. In section 3 we study the generation of the flavored $B - L_i$ asymmetries showing how the inclusion of light flavors in the analysis may impact the final asymmetry. Section 3.1 is devoted to the analysis of high scale (T_2) leptogenesis, section 3.2 to the possible collider patterns of T_1 , induced by requiring that the asymmetry generated at a higher scale survive T_1 related washouts, and in section 3.3 we discuss the interplay between conditions for successful high scale leptogenesis and lepton flavor violating observables present in minimal type III see-saw models. In section 4 we summarize and present our conclusions. In appendix A we present our conventions and definitions used in our numerical calculations.

2 Generalities

From a bottom-up approach the type III see-saw model is a simple extension of the SM that contains N_T additional fermionic $SU(2)$ triplets (T_α) with vanishing hypercharge ¹. In the basis in which the Majorana mass matrix for the triplets M_{T_α} is real and diagonal the interactions induced by the new states are given by the following Lagrangian

$$\mathcal{L}_T = i \text{Tr} \bar{T}_\alpha \not{D} T_\alpha - \lambda_{i\alpha}^* \bar{\ell}_i T_\alpha \tilde{H} - \frac{1}{4} \text{Tr} T_\alpha^\dagger C M_{T_\alpha} T_\alpha + \text{h.c.} , \quad (1)$$

¹Consistency with neutrino data [25] requires at least two triplets. However, apart from this constraint, the number of fermion triplets is arbitrary.

where $\ell = (\nu, l)^T$ and $H = (h^+, h^0)^T$ are the lepton and Higgs $SU(2)$ doublets ($\tilde{H} = i\tau_2 H^*$). The triplets can be written as a matrix

$$T_\alpha = \tau^A T_\alpha^A = \begin{pmatrix} T_\alpha^0 & \sqrt{2}T_\alpha^+ \\ \sqrt{2}T_\alpha^- & -T_\alpha^0 \end{pmatrix}, \quad (2)$$

where $T^0 = T^3$, $T^\pm = (T^1 \mp iT^2)/\sqrt{2}$ and C is the charge conjugation operator. In this notation, the covariant derivative is defined as $D_\mu = \partial_\mu - ig\tau^A W_\mu^A/2$. The Greek indices $\alpha, \beta \dots = 1, \dots, N_T$ are used to label the different fermionic triplet generations, Latin indices i, j, \dots for the lepton flavors e, μ, τ .

The Majorana triplets' mass terms break lepton number and thus, after electroweak symmetry breaking, induce Majorana neutrino masses for the left-handed neutrinos. In the basis (ν_{Li}, T_α^0) , where the heavy neutrino mass matrix is diagonal, the neutral $(3 + N_T) \times (3 + N_T)$ fermion mass matrix reads

$$\mathbf{M}_N = \begin{pmatrix} 0 & v \boldsymbol{\lambda} \\ v \boldsymbol{\lambda}^T & \hat{\mathbf{M}}_T \end{pmatrix}, \quad (3)$$

where v is the SM Higgs vacuum expectation value, $v \simeq 174$ GeV. Accordingly the effective light neutrino mass matrix is given by

$$\mathbf{m}_\nu^{eff} = -v^2 \boldsymbol{\lambda} \cdot \hat{\mathbf{M}}_T^{-1} \cdot \boldsymbol{\lambda}^T. \quad (4)$$

In general the Yukawa coupling matrix $\boldsymbol{\lambda}$ is a complex $3 \times N_T$ matrix in flavor space and therefore contains new sources of CP violation, as can be clearly seen by expressing $\boldsymbol{\lambda}$ by means of the Casas-Ibarra parametrization [26]:

$$\boldsymbol{\lambda} = \mathbf{U}^* \cdot \sqrt{\hat{\mathbf{m}}_\nu} \cdot \mathbf{R} \cdot \sqrt{\hat{\mathbf{M}}_T}, \quad (5)$$

which ensures that the correct low energy parameters are obtained. Here \mathbf{U} is the leptonic mixing matrix which diagonalizes the effective neutrino mass matrix \mathbf{m}_ν^{eff} and is fixed by low energy observables (three light neutrino mixing angles, one Dirac and one or two Majorana CP violating phases). The matrix \mathbf{R} is an orthogonal complex matrix defined by $N_T(N_T - 1)/2$ complex parameters.

The CP violating sources contained in $\boldsymbol{\lambda}$ induce two-body CP violating decays of T_α . The tree level decay width for these processes reads

$$\Gamma_\alpha = \sum_j \Gamma(T_\alpha \rightarrow \bar{\ell}_j \tilde{H}, \ell_j \tilde{H}^\dagger) = \frac{M_{T_\alpha}}{8\pi} \sum_{i=e,\mu,\tau} \lambda_{i\alpha}^* \lambda_{i\alpha} = \frac{M_{T_\alpha}^2}{8\pi v^2} \sum_{i=e,\mu,\tau} \tilde{m}_{i\alpha}, \quad (6)$$

where

$$\tilde{m}_{i\alpha} = \frac{v^2}{M_{T_\alpha}} \lambda_{i\alpha}^* \lambda_{i\alpha} \quad \text{and} \quad \tilde{m}_\alpha = \sum_{i=e,\mu,\tau} \tilde{m}_{i\alpha}. \quad (7)$$

In terms of (5) these parameters are determined by low energy observables and by the entries of the matrix \mathbf{R} . Given that such entries can be arbitrarily large, without affecting

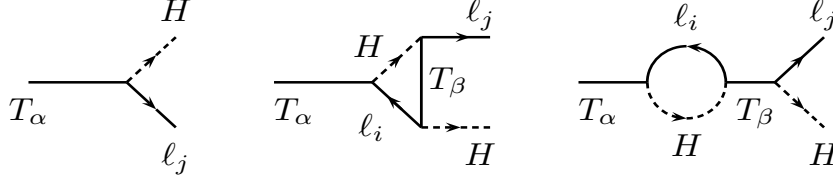


Figure 1: Diagrams generating the flavored CP violating asymmetries $\epsilon_{T_\alpha}^{\ell_i}$.

the light neutrino masses, light triplets ($\mathcal{O}(\text{TeV})$) not necessarily imply small values of \tilde{m}_α .

The CP violating asymmetries for T_α decays into ℓ_j lepton flavor arise, as in type I see-saw [27], through the interference between the tree level and the one-loop vertex (V) and wave-function (W) diagrams (see fig. 1). The corresponding expressions are given by

$$\begin{aligned}\epsilon_{T_\alpha}^{\ell_j(\text{V})} &= \frac{1}{8\pi} \sum_{\beta \neq \alpha} \frac{\text{Im}[\sqrt{\omega_\beta}(\lambda^\dagger \lambda)_{\beta\alpha} \lambda_{j\beta}^* \lambda_{j\alpha}]}{(\lambda^\dagger \lambda)_{\alpha\alpha}} f(\omega_\beta), \\ \epsilon_{T_\alpha}^{\ell_j(\text{W})} &= -\frac{1}{8\pi} \sum_{\beta \neq \alpha} \frac{\text{Im}\{[(\lambda^\dagger \lambda)_{\alpha\beta} + \sqrt{\omega_\beta}(\lambda^\dagger \lambda)_{\beta\alpha}] \lambda_{j\beta}^* \lambda_{j\alpha}\}}{(\lambda^\dagger \lambda)_{\alpha\alpha}} g(\omega_\beta),\end{aligned}\quad (8)$$

where $\omega_\beta = M_{T_\beta}^2/M_{T_\alpha}^2$ and the loop functions are given by

$$\begin{aligned}f(\omega_\beta) &= (1 + \omega_\beta) \ln\left(\frac{\omega_\beta + 1}{\omega_\beta}\right) - 1, \\ g(\omega_\beta) &= \frac{\omega_\beta - 1}{(\omega_\beta - 1)^2 + a_{\Gamma_\beta}^2},\end{aligned}\quad (9)$$

with $a_{\Gamma_\beta} = \Gamma_\beta/M_\alpha$. In the case of a quasi-degenerate triplet spectrum ($\sqrt{\omega_\beta} \sim 1 + a_{\Gamma_\beta}$) the wave-function piece is resonantly enhanced. In such cases, depending on the size of the CP violating phases, the flavored CP asymmetries, entirely determined by the wave-function piece, can be $\mathcal{O}(1)$ (see ref. [23, 24] for a thorough discussion of such possibility within type I see-saw). In the case of a hierarchical triplet spectrum ($\omega_\beta \gg 1$ with $\beta > \alpha$) a_{Γ_β} can be neglected and the loop functions expanded in powers of ω_β^{-1} . At leading order the flavored CP violating asymmetries can be expressed as

$$\epsilon_{T_\alpha}^{\ell_j} = -\frac{1}{8\pi(\lambda^\dagger \lambda)_{\alpha\alpha}} \sum_{\beta \neq \alpha} \text{Im} \left\{ \left[\frac{(\lambda^\dagger \lambda)_{\alpha\beta}}{\omega_\beta} + \frac{(\lambda^\dagger \lambda)_{\beta\alpha}}{2\sqrt{\omega_\beta}} \right] \lambda_{j\beta}^* \lambda_{j\alpha} \right\}. \quad (10)$$

Some comments are in order regarding the different terms in this expression. Since the term $\sum_\beta (\lambda^\dagger \lambda)_{\alpha\beta} (\lambda^\dagger \lambda)_{\beta\alpha}$ is real, the first term only contributes to the flavored CP violating asymmetries [11, 27]. Although, apart from the kinetic term, the Lagrangian in eq. (1) has the same structure as the Lagrangian in type I see-saw, the contractions of the $SU(2)$ indices in the Yukawa interaction terms are different and consequently the second term in (10) is a factor 3 smaller than the corresponding one in the standard case².

²In contrast to the type I see-saw case, due to these contractions, there is a relative sign between the wave-function and vertex pieces. Thus, at leading order, $\epsilon_{T_\alpha}^{\ell_j(\text{W})} \rightarrow 1/\sqrt{\omega_\beta}$ whereas $\epsilon_{T_\alpha}^{\ell_j(\text{V})} \rightarrow -1/2\sqrt{\omega_\beta}$

In the strongly hierarchical limit, where the first term in (10) can be neglected, and assuming a normal hierarchical light neutrino spectrum it can be shown that as in the standard case there is an upper bound on the flavored CP violating asymmetries, namely [10]

$$\epsilon_{T_\alpha}^{\ell_j} \lesssim 10^{-5} \left(\frac{M_{T_\alpha}}{10^{10} \text{ GeV}} \right) \left(\frac{m_3}{1 \text{ eV}} \right) \frac{\tilde{m}_{j\alpha}}{\tilde{m}_\alpha}. \quad (11)$$

Finally, the total CP violating asymmetry can be obtained from (10) by summing over the flavor index j [21]

$$\epsilon_{T_\alpha} = -\frac{1}{16\pi} \sum_\beta \frac{1}{\sqrt{\omega_\beta}} \frac{\text{Im}[(\lambda^\dagger \lambda)_{\beta\alpha}^2]}{(\lambda^\dagger \lambda)_{\alpha\alpha}}. \quad (12)$$

The major quantitative difference between the standard leptogenesis scenario and fermion triplet leptogenesis lies in the fact that fermion triplets couple to SM gauge bosons. At $z = M_T/T \ll 1$ gauge reactions thermalize the triplet distribution meaning that no asymmetry is produced at high T . Gauge reactions decouple at temperatures when

$$\frac{\Gamma_A}{H} = \frac{\gamma_A}{n_T^{\text{Eq}} H} \lesssim 1, \quad (13)$$

where γ_A is the gauge reaction density, H is the expansion rate of the Universe and n_T^{Eq} is the equilibrium triplet number density (see appendix A). Thus, if at this stage inverse decay processes $\ell H \rightarrow T$ are decoupled as well ($\gamma_D/n_\ell^{\text{Eq}} H \lesssim 1$, where γ_D is the decay reaction density), the CP violating out-of-equilibrium decays of the triplets will produce a $B - L$ asymmetry. Conversely, if inverse decays are still active when gauge reactions decouple the $B - L$ asymmetry will be generated at lower temperatures, after inverse decays switch off. Since $\gamma_A/\gamma_D \sim g^4/M_T \tilde{m}$, as M_T decreases only large values of \tilde{m} are able to maintain inverse decays active after gauge reaction decoupling takes place. Fig. 2 shows an illustrative example where it can be seen that for $M_T = 10^{12}$ GeV, inverse decay processes are still active after gauge interaction decoupling provided $\tilde{m} \sim 5 \times 10^{-3}$ eV, whereas for $M_T = 10^{10}$ GeV this occurs for a \tilde{m} an order of magnitude larger. Independently of how leptogenesis proceeds, a precise determination of the generated $B - L$ asymmetry requires a numerical treatment of the corresponding Boltzmann equations, which we now discuss in turn.

2.1 Generation of the $B - L$ asymmetry

Although in this section we will focus on the case in which the asymmetry is generated along one specific flavor direction, we will write flavor dependent Boltzmann equations as they will be used in sec. 3. Considering only triplet annihilations (mediated by SM gauge bosons), decays (induced by the Yukawa couplings λ) and $\Delta L = 2$ scatterings ($\ell_j \tilde{H}^\dagger \leftrightarrow \ell_i \tilde{H}^\dagger$ and $\bar{\ell}_j \tilde{H} \leftrightarrow \bar{\ell}_i \tilde{H}$)³ the network of flavor dependent Boltzmann equations

³Although we work at $\mathcal{O}(\lambda^2)$ the inclusion of these processes is mandatory to obtain Boltzmann equations with the correct thermodynamical behavior [28].

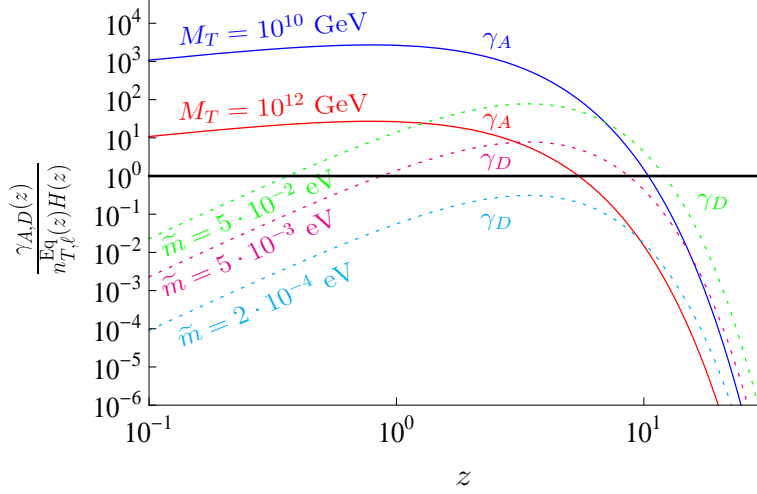


Figure 2: Evolution of gauge and Yukawa reaction densities (γ_A , γ_D) with z . The horizontal black line indicates when do either gauge or Yukawa reactions go out of equilibrium.

at $\mathcal{O}(\lambda^2)$, assuming a vanishing Higgs number asymmetry [12], can be written as

$$\frac{dY_{T_\alpha}}{dz_\alpha} = -\frac{1}{sH z_\alpha} \left[\left(\frac{Y_{T_\alpha}}{Y_{T_\alpha}^{\text{Eq}}} - 1 \right) \gamma_{D_\alpha} + \left(\frac{Y_{T_\alpha}^2}{(Y_{T_\alpha}^{\text{Eq}})^2} - 1 \right) \gamma_{A_\alpha} \right], \quad (14)$$

$$\frac{dY_{\Delta_i}}{dz_\alpha} = -\frac{1}{sH z_\alpha} \left[\left(\frac{Y_{T_\alpha}}{Y_{T_\alpha}^{\text{Eq}}} - 1 \right) \epsilon_{T_\alpha}^{\ell_i} + \frac{K_{i\alpha}}{2Y_{T_\alpha}^{\text{Eq}}} \sum_{j=e,\mu,\tau} C_{ij}^{\ell_i} Y_{\Delta_j} \right] \gamma_{D_\alpha}. \quad (15)$$

Here $z_\alpha = M_{T_\alpha}/T$, $Y_X = n_X/s$ and $Y_{\Delta_i} = Y_{\Delta_{B/3-L_i}}$ with $Y_{L_i} = 2Y_{\ell_i} + Y_{e_i}$ (the lepton asymmetry distributed in left and right handed degrees of freedom). The flavor projectors $K_{i\alpha}$ are defined as follows [12]:

$$K_{i\alpha} = \frac{\lambda_{i\alpha}^* \lambda_{i\alpha}}{(\lambda^\dagger \lambda)_{\alpha\alpha}} = \frac{\tilde{m}_{i\alpha}}{\tilde{m}_\alpha}, \quad (16)$$

note that $\sum_{i=e,\mu,\tau} K_{i\alpha} = 1$. The numerical coefficients $C_{ij}^{\ell_i}$, that relate Y_{ℓ_i} with Y_{Δ_i} according to [11]

$$Y_{\ell_i} = -\sum_j C_{ij}^{\ell_i} Y_{\Delta_j}, \quad (17)$$

couple the different differential equations in (15) and are determined by the reactions that at a certain temperature regime are in equilibrium [8, 11, 10], finally

$$\gamma_{D_\alpha} = \sum_{i=e,\mu,\tau} \gamma_{D_{i\alpha}}. \quad (18)$$

By integrating eq. (15) the resulting $Y_{\Delta_{B-L}}$ asymmetry can be expressed according to (see appendix A)

$$Y_{\Delta_{B-L}} = \sum_{i=e,\mu,\tau} Y_{\Delta_i} = 3 \times \sum_{i=e,\mu,\tau} \epsilon_{T_\alpha}^{\ell_i} Y_{T_\alpha}^{\text{Eq}} \eta_{i\alpha}, \quad (19)$$

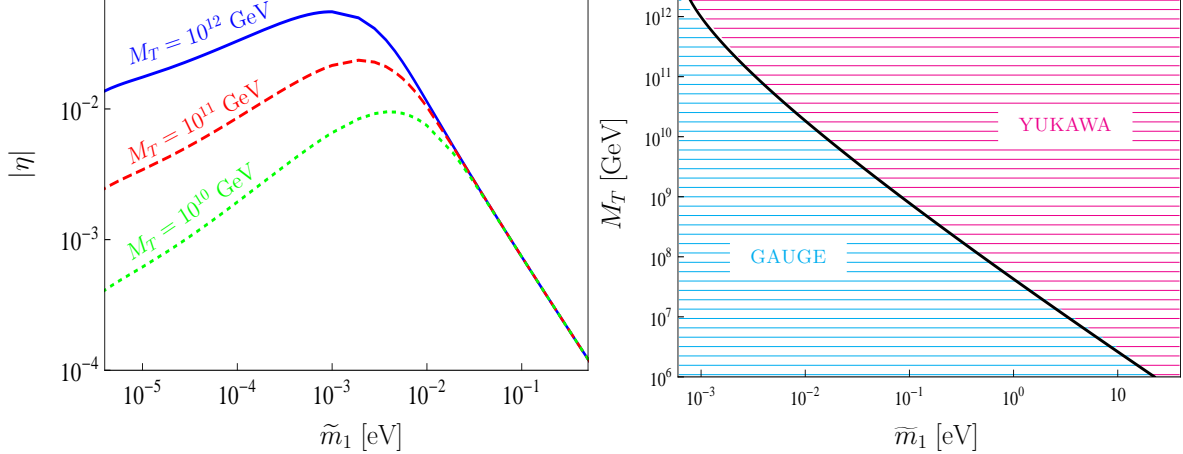


Figure 3: Efficiency factor as a function of \tilde{m}_1 in the flavor aligned case (left panel) and regions for which gauge interactions freeze out after (lower region) and before (upper region) Yukawa reaction decoupling (right panel).

where the factor 3 accounts for the three $SU(2)$ degrees of freedom of the triplets and $\eta_{i\alpha}$ denotes the efficiency factor that accounts for the T_α generated asymmetry in flavor i that survives washouts. In the case of a hierarchical spectrum ($M_{T_1} \ll M_{T_{\alpha>1}}$), neglecting the dynamics of the heavier triplets, leptogenesis proceeds entirely through T_1 dynamics. Compared with the standard case the amount of $Y_{\Delta_{B-L}}$ is smaller and its value can strongly depend on M_{T_1} (depending on the size of \tilde{m}_1). In order to demonstrate this we have numerically solved eqs. (14) and (15) assuming alignment along flavor j ($K_{j1} = \delta_{j1}$) and with $C_{ij}^\ell = \delta_{ij}$. Figure 3 (left panel) shows the evolution of the efficiency factor (η) with \tilde{m}_1 for several values of M_{T_1} . It can be seen that the dependence on M_{T_1} is stronger for small values of \tilde{m}_1 and less pronounced as \tilde{m}_1 increases, being completely independent of M_{T_1} for large values of \tilde{m}_1 , which means there is a \tilde{m}_1^{\min} above which leptogenesis proceeds as in the standard case⁴. This implies that although gauge processes introduce a dependence on the triplet mass this dependence disappears as soon as the gauge interactions are subdominant with respect to Yukawa reactions. Thus, in the fermionic triplet leptogenesis the generation of the $B - L$ asymmetry can proceed either in a region determined by the condition $\tilde{m}_1 < \tilde{m}_1^{\min}$ or conversely in a region defined by $\tilde{m}_1 > \tilde{m}_1^{\min}$. The determination of \tilde{m}_1^{\min} , for a given triplet mass, can be done as follows. From the gauge decoupling condition $\gamma_A/n_T^{\text{Eq}}H \lesssim 1$ the corresponding $z = z_{A-\text{dec}}$ at which gauge reactions go out of equilibrium can be calculated. The value \tilde{m}_1^{\min} can be computed by requiring that the Yukawa interactions are still active at $z_{A-\text{dec}}$ i.e. $(\gamma_D/n_\ell^{\text{Eq}}H)|_{z=z_{A-\text{dec}}} \gtrsim 1$. Figure 3 shows the values (\tilde{m}_1, M_{T_1}) for which the $B - L$ asymmetry either depends on M_{T_1} (“gauge region”) or is entirely determined by \tilde{m}_1 (“Yukawa region”).

⁴In standard leptogenesis at $\mathcal{O}(\lambda^2)$ the efficiency does not depend on the fermionic singlet mass at all.

3 Including Flavor

Flavor effects become relevant at temperatures below 10^{13} GeV when bottom and tau Yukawa interactions enter into thermodynamic equilibrium [8, 10, 11]. For definiteness, from now on, we will focus in the temperature window 10^9 GeV $\lesssim M_{T_1} \lesssim 10^{12}$ GeV at which, in addition to the bottom and tau Yukawa processes, also electroweak sphalerons are already in thermal equilibrium. In this regime the $B - L$ asymmetry is distributed along ℓ_τ and ℓ_1 (an admixture of muon and electron flavors). The determination of the total asymmetry in this case is therefore a two flavor problem (ℓ_τ , ℓ_1) and the network of Boltzmann equations in (14)-(15) consist of three coupled differential equations in the variables Y_{T_1} , Y_{Δ_τ} and Y_{Δ_1} where the flavored asymmetries are coupled by the flavor coupling matrix [11]

$$C^\ell = \frac{1}{460} \begin{pmatrix} 196 & -24 \\ -9 & 156 \end{pmatrix}, \quad (20)$$

obtained from the chemical equilibrium conditions imposed by reactions that in the corresponding temperature window are faster than the expansion rate [29]. According to the discussion of the previous section, in triplet fermion leptogenesis the asymmetry can proceed in one of the two following *regimes*: (a) Gauge reaction decoupling takes place after Yukawa induced inverse decays have gone out of equilibrium – “gauge region”; (b) Gauge reactions freeze-out when inverse decay processes are still active – “Yukawa region”. In scenario (a) since the the triplet abundance is efficiently diminished by gauge boson mediated annihilations the effects of flavor are expected to be small, or even negligible. Instead, in scenario (b) since at $z_1 \gg 1$ the dynamics of T_1 is entirely determined by Yukawa reactions, flavor effects may be sizable and can be prominent as long as

$$\left| \epsilon_{T_1}^{\ell_j} \right| > \left| \epsilon_{T_1}^{\ell_i} \right| \quad \text{and} \quad K_{j1} \gg K_{i1}. \quad (21)$$

In order to show this is in fact the case, we have calculated the evolution of the $Y_{\Delta_{B-L}}(z)$ asymmetry for $M_{T_1} = 10^{12}$ GeV and $\tilde{m}_1 = 2 \times 10^{-4}, 5 \times 10^{-2}$ eV by numerically integrating eqs. (14) and (15), and assuming the following flavor configuration $K_{11} = 0.99$, ($K_{\tau 1} = 1 - K_{11}$) and $\epsilon_{T_1}^{\ell_1} = -0.1 \times \epsilon_{T_1}$ and $\epsilon_{T_1}^{\ell_\tau} = 1.1 \times \epsilon_{T_1}$ with $\epsilon_{T_1} = 10^{-5}$. The result is displayed in fig. 4 where, in addition to the resulting asymmetry in the flavored case, we also show the corresponding asymmetry in the fully aligned case to facilitate the comparison between both cases. As can be seen in fig. 4 (left panel) for $\tilde{m}_1 = 2 \times 10^{-4}$ eV (regime (a)) flavor effects are small, producing only a $\sim 5\%$ enhancement. In contrast, for $\tilde{m}_1 = 5 \times 10^{-2}$ eV, and given that the flavor configuration we chose satisfies (21), flavor effects produce a two orders of magnitude enhancement of the final $B - L$ asymmetry: $Y_{\Delta_{B-L}}^{\text{Alig}} = 2.1 \times 10^{-10}$ whereas $Y_{\Delta_{B-L}}^{\text{Flav}} = 3.1 \times 10^{-8}$. As in the standard case the effects of flavor could be even larger once the muon Yukawa coupling enters into thermodynamical equilibrium since in that case two flavor projectors can be simultaneously small.

As previously discussed, the minimum value of \tilde{m}_1 for which flavor effects can produce a sizable enhancement of the $B - L$ asymmetry is not unique and depends on the triplet mass. Fig. 5 shows the behavior of the $B - L$ asymmetry with \tilde{m}_1 . As can be seen for $M_T = 10^{12}$ GeV flavor effects start playing a role at $\tilde{m}_1 \sim 10^{-3}$ eV and produce a

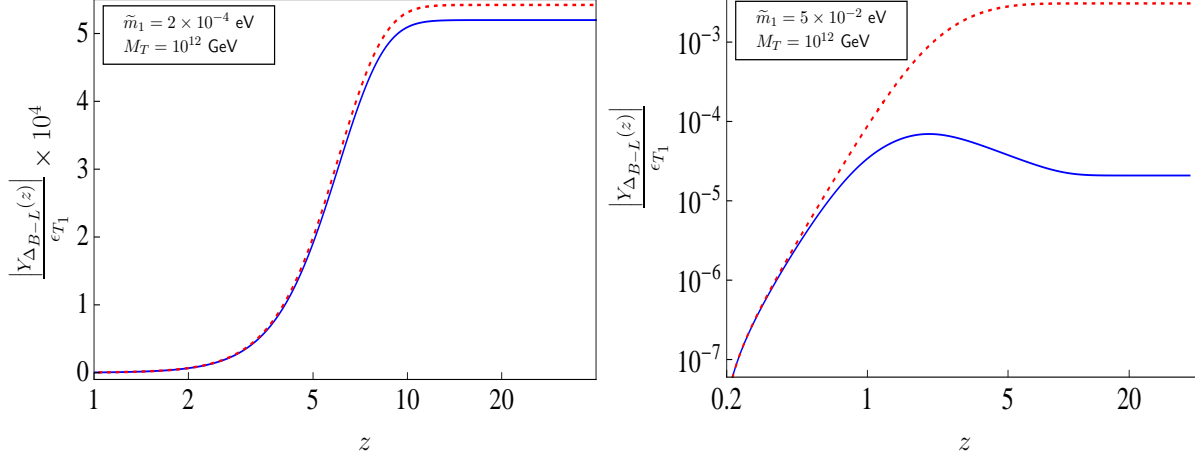


Figure 4: $Y_{\Delta_{B-L}}(z)/\epsilon_{T_\alpha}$ vs z for aligned (solid blue line) and flavored (dotted red line) cases in the temperature regime $10^{10} \text{ GeV} \lesssim T \lesssim 10^{12} \text{ GeV}$. The flavored CP asymmetries are fixed as $\epsilon_{T_1}^{\ell_1} = -0.1 \times \epsilon_{T_1}$, $\epsilon_{T_1}^{\ell_\tau} = 1.1 \times \epsilon_{T_1}$ and $\epsilon_{T_1} = 10^{-5}$ whereas the flavor projectors are accordingly $K_{11} = 0.99$ and $K_{1\tau} = 0.01$.

strong enhancement above this value i.e. well inside the “Yukawa region” (see fig. 3 right panel). For $M_T = 10^{10} \text{ GeV}$, as anticipated, the related effects start being relevant for larger values of \tilde{m}_1 , again well within the “Yukawa region”.

3.1 Implications for TeV leptogenesis

TeV triplets are thermalized by gauge boson mediated annihilations up to $z \gg 1$ (see fig. 2). The generation of the $B - L$ asymmetry in that case proceeds basically above this z once the relic fraction that survives annihilation start decaying. Sphaleron interactions transform this asymmetry into a B asymmetry up to temperatures T_{dec} at which their reactions are suddenly decoupled by the spontaneous breaking of the $SU(2)$ symmetry [30]. This constraint combined with $Y_{\Delta_B} \sim 10^{-10}$ implies the bound $M_T \gtrsim 1.6 \text{ TeV}$ [22]. This is to be compared with the standard resonant leptogenesis framework in which the fermionic singlet can have a mass well below the TeV scale. The reason is that while in the standard case the efficiency is basically determined by \tilde{m} , in the fermionic triplet scenario there is a dependence on M_T that strongly suppress the efficiency when $M_T \sim \mathcal{O}(\text{TeV})$.

Adding flavor effects in the analysis does not allow to relax the bound: in sec. 3 we argued that flavor effects can yield a relevant enhancement of the $B - L$ asymmetry (flavored efficiencies) only when gauge interactions decouple before Yukawa reactions, which requires large values of \tilde{m} , as shown in fig. 3 (right panel). In order to discuss how this could take place for TeV triplets let us fix $M_T = 10^3 \text{ GeV}$. Since the sphaleron decoupling temperature is determined by the Higgs mass (m_h) according to $T_{\text{dec}} \simeq [80 + 0.45(m_h/\text{GeV})] \text{ GeV}$ (with $m_h \in [114, 200] \text{ GeV}$) [22, 30], fixing $m_h = 120 \text{ GeV}$ sphaleron decoupling will take place at $z_{\text{dec}} \sim 7.5$. At this z Yukawa reactions will overcome gauge

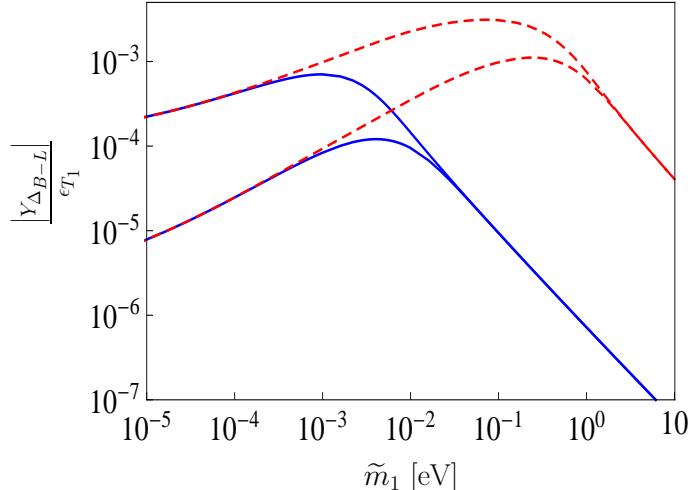


Figure 5: $Y_{\Delta_{B-L}}/\epsilon_{T_1}$ as a function of \tilde{m}_1 for flavored (dotted red lines) and aligned (solid blue line) cases. The upper (lower) dotted and solid curves correspond to $M_{T_1} = 10^{12}$ GeV ($M_{T_1} = 10^{10}$ GeV). The flavored CP asymmetries and flavor projectors were fixed as in fig. 4.

interactions provided $\tilde{m} \gtrsim 3.5$ keV⁵, however since no asymmetry could be produced until inverse Yukawa decays decouple flavor effects are irrelevant, as this will take place much more above z_{dec} .

Given that the LHC reach for fermion triplets discovery is around 700-800 GeV [31, 32], from the previous discussion, one could be tempted to conclude that any observation of fermionic triplets at the LHC would rule out leptogenesis within this framework. This however is not entirely correct as the asymmetry might be generated by degrees of freedom at higher scales – T_2 in triplet scenario for example⁶. In this case, since T_1 will not play any role in the generation of the $B - L$ asymmetry, the lower limit on M_{T_1} will no longer hold. However, although T_1 does not participate in leptogenesis, the asymmetry might still be erased by T_1 flavor dynamics. The particular flavor structures required to prevent too large T_1 related washouts may yield particular collider signals that can be used to determine whether leptogenesis might be relevant for the generation of the cosmic baryon asymmetry. In the following we will discuss such a possibility.

Without the loss of generality, we assume the asymmetry is produced by heavy fermion triplet (T_2) dynamics at scales 10^9 GeV $\lesssim M_{T_2} \lesssim 10^{12}$ GeV. Requiring $M_{T_1} \lesssim 1$ TeV the flavored CP violating asymmetries $\epsilon_{T_1}^{\ell_j}$ are negligible (see eq. (11)) and accordingly T_1 decays do not generate any asymmetry. Thus, given this triplet mass spectrum, the asymmetry is generated by T_2 decays (at $z_2 = M_2/T \sim 1$) in the two-flavored regime (ℓ_τ, ℓ_2), where ℓ_2 is an admixture of muon and electron flavors. The calculation of the $B - L$ asymmetry at this stage goes along the same lines as in the case we discussed at the beginning of this section, but taking into account that in this case we deal with

⁵Such values are possible only if the different Yukawa couplings λ are highly fine-tuned.

⁶This possibility has been studied in the standard case in both, the one flavor approximation [13] and flavored [14, 15] cases.

T_2 dynamics. At $z \ll z_2$, once muon Yukawa interactions enter in thermodynamical equilibrium, the Y_{Δ_τ} produced at $z \sim z_2$ remains frozen whereas Y_{Δ_2} splits into electron and muon asymmetries. The splitting, in the case of $Y_{T_2}^{\text{in}} = 0$, is determined by [15]

$$Y_{\Delta_k} = \frac{K_{k2}}{\widetilde{K}_{22}} Y_{\Delta_2}, \quad (22)$$

with $k = e, \mu$ and $\sum_{k=e,\mu} K_{k2}/\widetilde{K}_{22} = 1$.

As the temperature drops and reaches $z_1 = M_{T_1}/T \sim 1$, T_1 related washout effects become effective and are determined by

$$\frac{dY_{\Delta_i}(z_1)}{dz_1} = -\frac{K_{i1}}{2sH z_1 Y_\ell^{\text{Eq}}} \sum_{j=e,\mu,\tau} C_{ij}^\ell Y_{\Delta_j}(z_1) \gamma_{D_1} = -\frac{\kappa_{i1}}{4} \sum_{j=e,\mu,\tau} C_{ij}^\ell Y_{\Delta_j}(z_1) K_1(z_1) z_1^3, \quad (23)$$

where $K_1(z_1)$ is the modified Bessel function of first-type, the parameter $\kappa_{i1} = \widetilde{m}_{i1}/m_\star$, with $m_\star = 8\pi v^2 H|_{z_1=1}/M_{T_1}^2 = 2.25 \times 10^{-3}$ eV, and the flavor coupling matrix at this stage is given by [11]

$$C^\ell = \frac{1}{711} \begin{pmatrix} 221 & -16 & -16 \\ -16 & 221 & -16 \\ -16 & -16 & 221 \end{pmatrix}. \quad (24)$$

A rough estimate of washout effects can be done by simply neglecting flavor coupling i.e. setting $C^\ell = \mathbf{I}$ and assuming the T_1 related washout processes are efficient before sphalerons decouple below the electroweak phase transition. In that case eq. (23) can be analytically integrated and the final baryon asymmetry can be written in terms of flavored asymmetries ($Y_{\Delta_i}^{\text{in}}$) generated at a high scale (in T_2 decays) and T_1 washout related damping factors

$$Y_{\Delta_B} \simeq \sum_{i=e,\mu,\tau} Y_{\Delta_i}^{\text{in}} e^{-3\pi\kappa_{i1}/8}. \quad (25)$$

Whether the correct amount of baryon asymmetry can be generated therefore depends not only on $Y_{\Delta_i}^{\text{in}}$ but also on the T_1 related washout parameters κ_{i1} . In particular, for any value of $Y_{\Delta_i}^{\text{in}}$, there exist $\kappa_{i1} > \kappa_{i1}^{\text{max}}$, for which that particular flavored asymmetry will be completely washed out by T_1 dynamics before sphaleron decoupling and will thus not contribute to Y_{Δ_B} . Including flavor dynamics and finite sphaleron decoupling temperature, the values of κ_{i1}^{max} must be determined numerically, by solving eqs. (23) for a given set of $Y_{\Delta_i}^{\text{in}}$ and M_{T_1} . Generically however, taking into account that T_1 related washouts are relevant only if $\kappa_{i1} \gtrsim 1$ we can already distinguish between three cases:

1. **$\kappa_{i1} \ll 1$ for all flavors.** In that case all T_1 washout processes are weak and the baryon asymmetry is determined by high scale dynamics (of T_2 in our scenario). Successful leptogenesis is possible in principle depending on the (practically unmeasurable) high-scale (T_2 related) parameters.
2. **$\kappa_{i1} \ll 1$ and $\kappa_{j1} \gtrsim 1$.** The asymmetry in flavor(s) j will now be subject to possibly strong washouts due to T_1 . However, the asymmetry in flavor(s) i will again be completely determined at the high scale (by T_2 decays). In general therefore any

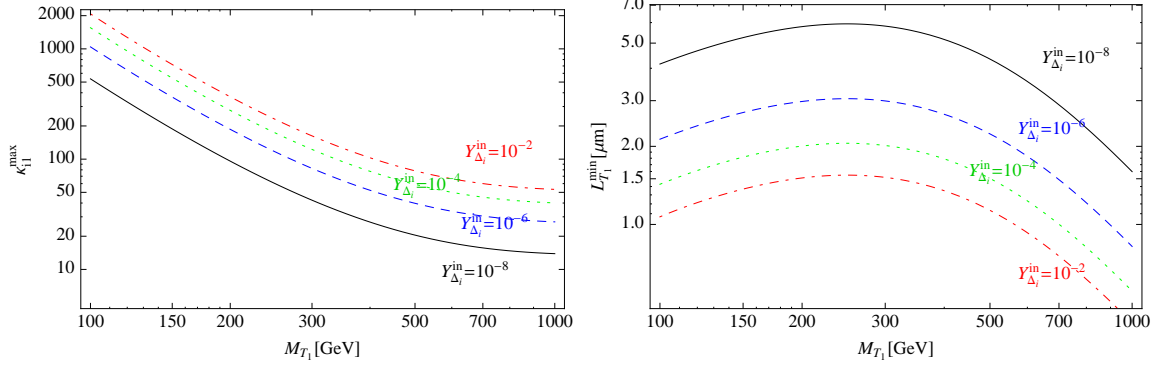


Figure 6: Solutions of κ_{i1}^{\max} and the minimum total T_1 decay length versus the T_1 mass for given values of high-scale generated lepton asymmetries $Y_{\Delta_i}^{\text{in}}$. See text for details.

constraints on κ_{j1} rely on the assumptions about the size of $Y_{\Delta_i}^{\text{in}}$. Even if some of $Y_{\Delta_j}^{\text{in}}$ would be strongly damped, leptogenesis might (or not) account for the baryon asymmetry.

3. $\kappa_{i1} \gtrsim 1$ for all flavors. Washout effects are relevant in all the flavors and the relevant washout parameters can in principle be constrained, depending on the size of $Y_{\Delta_i}^{\text{in}}$. Once for a given flavor κ_{i1} reaches κ_{i1}^{\max} the asymmetry in such a flavor will not contribute to Y_{Δ_B} implying that the final baryon asymmetry can be given by

- (a) Single flavor: $\kappa_{i1} < \kappa_{i1}^{\max}$ and $\kappa_{(j,k)1} > \kappa_{(j,k)1}^{\max}$
- (b) Two flavors: $\kappa_{(i,j)1} < \kappa_{(i,j)1}^{\max}$ and $\kappa_{k1} > \kappa_{k1}^{\max}$
- (c) Three flavors: $\kappa_{(e,\mu,\tau)1} < \kappa_{(e,\mu,\tau)1}^{\max}$

Finally, in order to determine κ_{i1}^{\max} we solve numerically eqs. (23) for given individual $Y_{\Delta_i}^{\text{in}}$ ranging from zero up to 10^{-2} (the largest values are motivated by completeness in order to cover possible extreme cases such as resonant leptogenesis and possibly large flavor enhancements) and require the final asymmetry $Y_{\Delta_B}(T_{\text{dec.}})$ to satisfy the experimental constraint $Y_{\Delta_B} = [8.52, 8.98] \times 10^{-11}$. We plot the obtained solutions (which are independent of lepton flavor, as can be inferred directly from the lepton flavor permutation symmetry of eq. (23)) as contours in the plane of M_{T_1} and κ_{i1}^{\max} on the left hand side of figure 6. We see that successful high scale thermal leptogenesis ($Y_{\Delta_i}^{\text{in}} \lesssim 10^{-8}$) requires $\kappa_{i1} < 10(500)$ for the T_1 masses of $M_{T_1} = 1000(100)$ GeV. On the other hand, in the extreme case $Y_{\Delta_i}^{\text{in}} \sim 10^{-2}$, T_1 washout parameters as large as $\kappa_{1i} \lesssim 50(2000)$ are allowed for T_1 masses of $M_{T_1} = 1000(100)$ GeV respectively.

3.2 Implications for type III see-saw models at the LHC

Experimentally some conclusions about whether successful high scale (due to T_2 or any other mechanism, possibly unrelated to neutrino mass generation) leptogenesis is achieved can be drawn by producing T_1 at colliders, verifying its lepton number violating interactions (typically using same sign di-lepton event signatures [31, 32]) and measuring its

branching ratios and the total decay width. From eq. (6) the triplet partial decay width can be written according to

$$\Gamma(\ell_i) \simeq 10^{-12} \kappa_i \left(\frac{M_{T_1}}{1 \text{ TeV}} \right)^2 \text{ GeV}. \quad (26)$$

In case 1 the triplet lifetime and the associated proper decay length in collider experiments is large, $L_{T_1} \gtrsim 50 (1 \text{ TeV}/M_{T_1})^2 \mu\text{m}$. Thus, an observation of a long lived triplet at the LHC can be fully consistent with high scale leptogenesis accounting for the Cosmic baryon asymmetry.

Excluding case 3c, cases 2 and 3 are conversely characterized by a short lived triplet and at least one (moderately) suppressed decay channel. Cases 3a, 3b and 3c are particularly interesting. First of all, the upper bound $\kappa_{(e,\mu,\tau)1} \lesssim \kappa_{i1}^{\text{max}}$ when all the flavored asymmetries contribute to Y_{Δ_B} implies a lower limit on the total T_1 proper decay length. The bound depends on M_{T_1} and is shown on the right hand side of figure 6. We observe that for $Y_{\Delta_i}^{\text{in}} \lesssim 10^{-8}$ (as allowed by high scale thermal non-resonant leptogenesis) the T_1 proper decay length should be larger than about $L_{T_1}^{\text{min}} \simeq 5 \mu\text{m}$ for the interesting range of M_{T_1} . This is close to the ATLAS inner tracker intrinsic accuracy [33], meaning that this region could possibly be identified by searching for displaced secondary vertices in the detector⁷. The prospects are somewhat thinner for the extreme cases of $Y_{\Delta_i}^{\text{in}} \lesssim 10^{-2}$, where the T_1 proper decay length should be larger than about $L_{T_1}^{\text{min}} \simeq 1 \mu\text{m}$ for the relevant M_{T_1} range.

Given the restrictions defining cases 3a and 3b it turns out that if one increases a single $\kappa_{i1} \gg \kappa_{i1}^{\text{max}}$, thus proportionally increasing the total triplet width, it follows that at least one flavored branching ratio needs to decrease and that the asymmetry will be produced mostly in that flavor. Experimentally the least favorable case is in the transition region between these two regimes, when the proper decay length of the triplets is still below the experimental resolution while the hierarchy between the different flavored branching ratios is the least, making definite statements about the viability of high scale leptogenesis difficult. This situation is visualized in fig. 7 for representative values of $Y_{\Delta_i}^{\text{in}}$ and for $M_{T_1} = 300 \text{ GeV}$. Note also that in order to discriminate between the cases 3a and 3b, and the parameter region where high scale leptogenesis is ruled out, all three leptonic decay channels of T_1 would need to be measured or at least experimentally constrained.

3.3 Implications for lepton flavor violation in minimal type III (+I) see-saw models

Recently, it has been shown, that a mixed type I+III see-saw model with a single fermionic singlet and triplet below the TeV scale naturally arises in a particular grand unified scenario [34, 35]. It belongs to a class of so-called ‘minimal’ type I and III see-saw scenarios, where the light neutrino mass matrix due to the mixing with only two heavy Majorana states is of rank 2 and the lightest neutrino remains (almost) massless. A particularity of such models is that the neutrino mass scale is fixed and that all the

⁷Note that in the detector reference frame, decay lengths of particles of a couple of hundred GeV may receive sizable boost factors at the LHC.

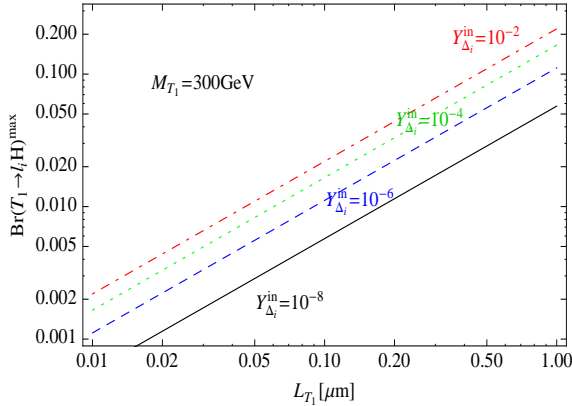


Figure 7: Parameter space regions of T_1 , allowing for successful high scale leptogenesis yielding the given asymmetry values in lepton flavor i . Regions below the curves are allowed for each $Y_{\Delta_i}^{\text{in}}$ value. $Br(T_1 \rightarrow \ell_i H)^{\text{max}}$ stands for the maximum allowed branching ratio of T_1 to a lepton of flavor i . See text for details.

interactions of the heavy mediators are fixed in terms of the measured neutrino oscillation data, a single physical Majorana phase and a complex parameter z [36], which is the single complex rotation angle in the 2×2 orthogonal matrix \mathbf{R} introduced in eq. (5) (we employ conventions as defined in [37]). The presence of TeV scale triplets can induce lepton flavor violating (LFV) processes which, in the minimal models, are all correlated and their magnitude scales exponentially with $\text{Im } z$. The greatest sensitivity is exhibited in the $\mu - e$ sector by the nuclear conversion experiments, yielding present bounds of the order $\text{Im } z < 7$ for the lightest triplet mass of $M_T = 100$ GeV [37]. The proposed next generation PRISM/PRIME experiments [38] are expected to improve the sensitivity down to $\text{Im } z < 4$ for the same triplet mass.

From the discussion in the previous sections it is clear, that minimal type III models at the TeV scale cannot mediate successful leptogenesis⁸. However, the requirements that the flavor dynamics of the TeV scale triplets do not wash out any lepton number asymmetries generated at a higher scale by other mechanisms can in these models be phrased entirely in terms of M_T and z . As such they can then be correlated with the projected sensitivity of future LFV experiments. In particular for the minimal I+III see-saw scenario we plot the constraints imposed by allowing for high scale leptogenesis for a given $Y_{\Delta}^{\text{in}} \equiv \sum_i Y_{\Delta_i}^{\text{in}}$ on the parameters of the model ($\text{Im } z$ and M_T) and superimpose them onto the projected sensitivity of future $\mu - e$ conversion experiments. In this procedure we marginalize over unknown phases, individual $Y_{\Delta_i}^{\text{in}}$ (while keeping their sum fixed) and also $\text{Re } z$. The results are shown in figure 8 for normal (on the left side) and inverted (on the right side) neutrino

⁸It has been proposed that resonant leptogenesis could be viable in the TeV scale minimal I+III model [35]. At first glance, it seems like an interesting possibility, since the fermionic singlet mediator would not suffer from the gauge reaction related strong washouts plaguing TeV triplet leptogenesis. However, at temperatures above the electroweak phase transition the mixing between $SU(2)$ singlets and triplets is highly suppressed and a large resonant enhancement due to the interference with the wavefunction diagram in figure 1 cannot occur.

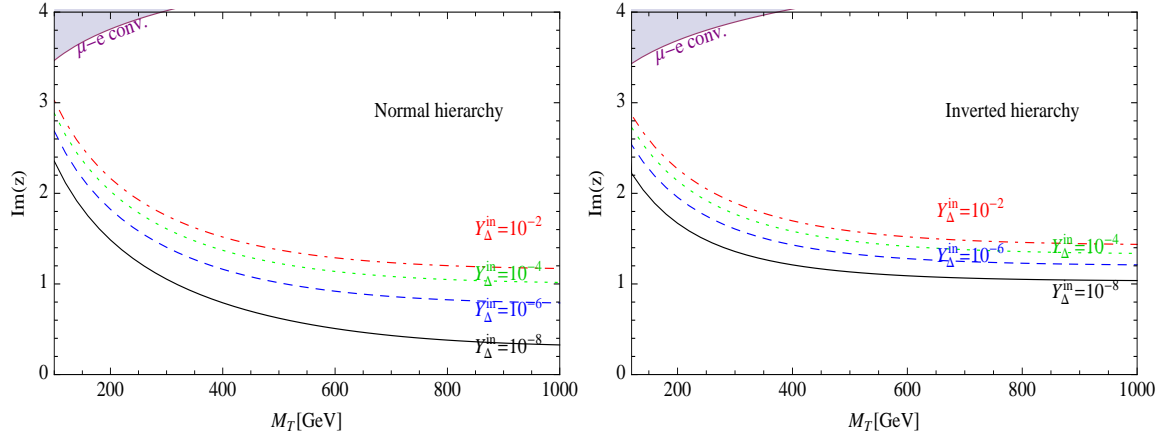


Figure 8: Correlation of leptogenesis bounds (area below the contours is allowed for a given Y_{Δ}^{in}) and projected $\mu - e$ conversion sensitivity of PRISM/PRIME [38] (shaded area) on the parameters of the minimal I+III model ($\text{Im}z$ and M_T). See text for details.

mass hierarchy. The results for the minimal pure type III models are similar. We observe that in the minimal models, requirements of successful high scale leptogenesis do not allow for LFV effects which could be detected in the foreseeable future. Conversely, observation of minimal see-saw models at the LHC and detection of the associated LFV signals in the next generation of $\mu - e$ nuclear conversion experiments would disfavor leptogenesis as the mechanism for the generation of the cosmic baryon asymmetry.

4 Conclusions

The generation of a lepton asymmetry, via the decays of fermionic $SU(2)$ triplets, can take place in two possible regimes: a “gauge regime” in which gauge reactions decouple at lower temperatures than Yukawa induced inverse decays. In that case the relic fraction of triplets that survives the gauge boson mediated annihilation processes will decay producing an amount of $B - L$ asymmetry whose value will strongly depend on M_T . Conversely, in the “Yukawa regime”, at low temperatures, the triplet distribution is driven to thermal equilibrium by Yukawa inverse decay processes, even after gauge interactions have switched off. In contrast to the “gauge regime”, in this case the amount of $B - L$ generated in T decays is not sensitive to M_T and is fixed by \tilde{m} . We have determined the $\tilde{m} - M_T$ regions that define both regimes. After identifying these regions we have studied the effects of light lepton flavors on the generation of the $B - L$ asymmetry. We have shown that in the “gauge regime” the effects have a minor influence as they barely reach fractions of a percent. Instead, in the “Yukawa regime”, as in the standard case, the enhancement of the final $B - L$ can be large, readily reaching the order of magnitude level.

We have shown that the requirement of successful T_1 leptogenesis constrains the dynamics of TeV triplets to lie in the “gauge regime”, and therefore the inclusion of light lepton flavor effects in the calculation of the $B - L$ asymmetry has basically no significant

effect. This implies that even after including flavor the bound $M_{T_1} \gtrsim 1.6$ TeV, found in the one flavor approximation [22], still holds. Since fermionic triplets will be produced at the LHC if their masses are below ~ 1 TeV [31] their observation will automatically preclude them from explaining the Cosmic baryon asymmetry via T_1 leptogenesis. However, it will not necessarily exclude the case for high scale leptogenesis, due to T_2 or other mechanisms. In fact, as we have argued, the $B - L$ asymmetry can be built up at a very early age ($T \gg M_{T_1}$) and can possibly survive T_1 washout processes. We have shown that T_1 collider observables could, in principle, be used either to strengthen or rule out the case for high scale leptogenesis⁹ (see also [39] for a recent related attempt to invalidate leptogenesis using collider observables). Finally, we have shown how in minimal see-saw models, LFV observables can also be employed in order to over-constrain the region of parameters of TeV scale Majorana electroweak triplets allowing for successful high scale leptogenesis.

5 Acknowledgements

We want to thank Enrico Nardi, Juan Racker and Ilja Doršner for many interesting conversations on the subject. Especially to Enrico Nardi for his suggestions and comments on the paper. D.A.S wants also to thank the theoretical physics group at the Josef Stefan Institute for their hospitality. D.A.S is supported by a Belgian FNRS postdoctoral fellowship. This work is supported in part by the Slovenian Research Agency, by the European Commission RTN network, Contract No. MRTN-CT-2006-035482 (FLAVIANet) and by the Deutsche Forschungsgemeinschaft via the Junior Research Group ‘‘SUSY Phenomenology’’ within the Collaborative Research Centre 676 ‘‘Particles, Strings and the Early Universe’’.

A Conventions and notation

In this appendix we fix our conventions and present the relevant equations used in the calculations discussed in sections 2 and 3. Some of the formulae are quite standard, however, for completeness, we include them. As has been mentioned we used Maxwell-Boltzmann distributions for the equilibrium number densities, namely

$$n_T^{\text{eq}}(z) = \frac{M_T^3}{\pi^2} \frac{K_2(z)}{z}, \quad n_\ell^{\text{eq}}(z) = \frac{2M_T^3}{\pi^2 z^3}, \quad (27)$$

where $z = M_T/T$ and $K_2(z)$ is the modified Bessel function of the second-type. With this approximation the energy density $\rho(z)$ and pressure $p(z)$ becomes

$$\rho(z) = \frac{3M_T^4}{z^4\pi^2} g_*, \quad p(z) = \frac{M_T^4}{z^4\pi^2} g_* \quad (28)$$

⁹Note that a similar argument could apply for other mechanisms of high scale baryogenesis, since sphaleron processes will tend to equilibrate any B asymmetry with L , which could in term be washed out by light Majorana triplet interactions.

where $g_* = \sum_{i=\text{All species}} g_i$ is the number of standard model relativistic degrees of freedom (118 for $T \gg 300$ GeV). Accordingly, the expansion rate of the Universe and entropy density can be written as

$$H(z) = \sqrt{\frac{8g_*}{\pi}} \frac{M_T^2}{M_{\text{Planck}}} \frac{1}{z^2}, \quad s(z) = \frac{4M_T^3}{z^3\pi^2} g_*. \quad (29)$$

The decay and annihilation reaction densities in eqs. (14) and (15) are given by

$$\gamma_{D_\alpha}(z) = \frac{1}{8\pi^3} \frac{M_{T_\alpha}^5}{v^2} \frac{K_1(z)}{z} \tilde{m}_\alpha \quad (30)$$

$$\gamma_{A_\alpha}(z) = \frac{M_{T_\alpha}^4}{64\pi^4} \int_4^\infty dx \frac{\sqrt{x} K_1(z\sqrt{x}) \hat{\sigma}_A(x)}{z}, \quad (31)$$

where $K_1(z)$ is the modified Bessel function of first-type and $x = s/M_T^2$. The reduced cross section in (31) accounts for the gauge boson mediated s-channel processes $T_\alpha T_\alpha \leftrightarrow \ell\bar{\ell}$ and $T_\alpha T_\alpha \leftrightarrow q\bar{q}$ and the t and u-channel triplet mediated process $T_\alpha T_\alpha \leftrightarrow A_\mu A^\mu$ ($A^\mu = W_A^\mu, B^\mu$). For $\hat{\sigma}_A(x)$ we used the expression given in ref. [21]:

$$\hat{\sigma}_A(x) = \frac{6g^4}{\pi} \left(1 + \frac{2}{x}\right) r + \frac{2g^4}{\pi} \left[3 \left(1 + \frac{4}{x} - \frac{4}{x^2}\right) \log\left(\frac{1+r}{1-r}\right) - \left(4 + \frac{17}{x}\right) r\right] \quad (32)$$

with $r = \sqrt{1 - 4/x}$.

Neglecting flavor coupling, i.e $C^\ell = I$ eq. (15) can be analytically integrated (by using its integrating factor). Assuming a vanishing initial asymmetry ($Y_{\Delta_i}^{\text{in}}(z_0) = 0$) the result reads

$$Y_{\Delta_i}(z) = -\epsilon_{T_\alpha}^{\ell_i} Y_{T_\alpha}^{\text{Eq}}(z_0) \eta_{i\alpha}(z). \quad (33)$$

The flavored efficiency function $\eta_{i\alpha}(z)$ determines the evolution of the Δ_i asymmetry at any z . Using the approximation $Y_{T_\alpha}(z) + Y_{T_\alpha}^{\text{Eq}}(z) \simeq 2Y_{T_\alpha}^{\text{Eq}}(z)$ it can be written as

$$\eta_{i\alpha}(z) = \frac{1}{Y_{T_\alpha}^{\text{Eq}}(z_0)} \int_{z_0}^z dz' \frac{\gamma_{D_\alpha}}{\gamma_{D_\alpha} + 2\gamma_{A_\alpha}} \frac{dY_{T_\alpha}(z')}{dz'} e^{-\sum_\alpha \int_{z'}^z dz'' P_{i\alpha}(z'')}, \quad (34)$$

where freeze-out of Δ_i takes place at z_f ($z_0 \ll z_f$) and

$$P_{i\alpha}(z) = \frac{K_{i\alpha}}{2Y_\ell^{\text{Eq}}(z)} \frac{\gamma_{D_\alpha}(z)}{s(z)H(z)z}. \quad (35)$$

For temperatures $T \gtrsim 10^{13}$ GeV, for which the flavor composition of the lepton doublet states is irrelevant, $K_{i\alpha} = 1$.

The final $B - L$ asymmetry is determined according to

$$Y_{\Delta_{B-L}} = \sum_{i=e,\mu,\tau} Y_{\Delta_i} = 3 \times \sum_{i=e,\mu,\tau} \epsilon_{T_\alpha}^{\ell_i} Y_{T_\alpha}^{\text{Eq}}(z_0) \eta_{i\alpha}(z_f), \quad (36)$$

where $\eta_{i\alpha}(z_f)$ defines the efficiency factor for flavor i and the factor 3 comes from the triplet degrees of freedom.

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