Implications of an additional scale on leptogenesis

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Abstract. We consider variations of the standard leptogenesis picture arising from the presence of an additional scale related to the breaking of a $U(1)_X$ abelian flavor symmetry. We show that quite generically the presence of an additional energy scale might introduce new qualitative and quantitative changes on leptogenesis. Especially interesting is the possibility of having succesful TeV leptogenesis with a vanishing total CP violating asymmetry. By solving the corresponding Boltzmann equations it is shown that these kind of scenarios encounters no difficulties in generating the Cosmic baryon asymmetry.

1. Introduction

From observations of light element abundances and of the Cosmic microwave background radiation [1] the Cosmic baryon asymmetry,

$$\mathcal{Y}_B = (8.75 \pm 0.23) \times 10^{-10}\,, \tag{1}$$

can be inferred. The conditions for a dynamical generation of this asymmetry (baryogenesis) are well known [2] and depending on how they are realized different scenarios for baryogenesis can be defined (see ref. [3] for a througout discussion).

Leptogenesis [4] is a scenario in which an initial lepton asymmetry, generated in the out-of-equilibrium decays of heavy standard model singlet Majorana neutrinos (N_α), is partially converted in a baryon asymmetry by anomalous sphaleron interactions [5] that are standard model processes. Singlet Majorana neutrinos are an essential ingredient for the generation of light neutrino masses through the seesaw mechanism [6]. This means that if the seesaw is the source of neutrino masses then qualitatively leptogenesis is unavoidable. Consequently, whether the baryon asymmetry puzzle can be solved within this framework turn out to be a quantitative question. This has triggered a great deal of interest on quantitative analysis of the standard leptogenesis model and indeed a lot of progress during the last years have been achieved (see ref. [7] for details).

Here we focus on variations of the standard leptogenesis picture which can arise if, apart from the lepton number breaking scale (M_N), an additional energy scale, related to the breaking of a new symmetry, exist. We consider a simple realization of this idea in which at an energy scale of the order of the lepton number violating scale the tree level coupling linking light and heavy neutrinos is forbiden by an exact $U(1)_X$ flavor symmetry which below (or above) M_N becomes spontaneously broken by the vacuum expectation value, σ, of a standard model singlet scalar

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field S and involves heavy vectorlike fields F_a. As will be discussed, according to the relative size of the relevant scales of the model (M_{N_α}, σ, M_{F_a}), different scenarios for leptogenesis can be defined [8]. Of particular interest is the case in which the total CP violating (CPV) asymmetries in the decays and scatterings of the singlet seesaw neutrinos vanish. As we will discuss further on two remarkable features distinguish this scenario [9]: (a) Flavor effects are entirely responsible for successful leptogenesis; (b) leptogenesis can be lowered down to the TeV scale.

The rest of this paper is organized as follows: In section 2 we briefly discuss the model while in section 3 we discuss two particular realizations of our scheme paying special attention to the *purely flavored leptogenesis* (PFL) case for which we analyse the evolution of the generated lepton asymmetry by solving the corresponding Boltzmann equations (BE).

2. The model

The model we consider here [8] is a simple extension of the standard model containing a set of $SU(2)_L \times U(1)_Y$ fermionic singlets, namely three right-handed neutrinos ($N_\alpha = N_{\alpha R} + N^c_{\alpha R}$) and three heavy vectorlike fields ($F_a = F_{aL} + F_{aR}$). In addition, we assume that at some high energy scale, taken to be of the order of the leptogenesis scale M_{N_1}, an exact $U(1)_X$ horizontal symmetry forbids direct couplings of the lepton ℓ_i and Higgs Φ doublets to the heavy Majorana neutrinos N_α. At lower energies, $U(1)_X$ gets spontaneously broken by the vacuum expectation value (vev) σ of a $SU(2)$ singlet scalar field S. Accordingly, the Yukawa interactions of the high energy Lagrangian read

$$-\mathcal{L}_Y = \frac{1}{2}\bar{N}_\alpha M_{N_\alpha} N_\alpha + \bar{F}_a M_{F_a} F_a + h_{ia}\bar{\ell}_i P_R F_a \Phi + \bar{N}_\alpha \left(\lambda_{\alpha a} + \lambda^{(5)}_{\alpha a}\gamma_5\right) F_a S + \text{h.c.} \qquad (2)$$

We use Greek indices $\alpha, \beta \cdots = 1,2,3$ to label the heavy Majorana neutrinos, Latin indices $a, b \cdots = 1,2,3$ for the vectorlike messengers, and $i, j, k, \ldots$ for the lepton flavors e, μ, τ. Following reference [8] we chose the simple $U(1)_X$ charge assignments $X(\ell_{L_i}, F_{L_a}, F_{R_a}) = +1$, $X(S) = -1$ and $X(N_\alpha, \Phi) = 0$. This assignment is sufficient to enforce the absence of $\bar{N}\ell\Phi$ terms, but clearly it does not constitute an attempt to reproduce the fermion mass pattern, and accordingly we will also avoid assigning specific charges to the right-handed leptons and quark fields that have no relevance for our analysis. The important point is that it is likely that any flavor symmetry (of the Froggatt-Nielsen type) will forbid the the same tree-level couplings, and will reproduce an overall model structure similar to the one we are assuming here. Therefore we believe that our results, that are focused on a new realization of the leptogenesis mechanism, can hint to a general possibility that could well occur also in a complete model of flavor.

As discussed in [8], depending on the hierarchy between the relevant scales of the model (M_{N_1}, M_{F_a}, σ), quite different *scenarios* for leptogenesis can arise. Here we will concentrate on two cases: (i) The standard leptogenesis case ($M_F, \sigma \gg M_N$); (ii) the PFL case ($\sigma < M_{N_1} < M_{F_a}$) that is, when the flavor symmetry $U(1)_X$ is still unbroken during the leptogenesis era and at the same time the messengers F_a are too heavy to be produced in N_1 decays and scatterings, and can be integrated away [9].

As is explicitly shown by the last term in eq. (2), in general the vectorlike fields can couple to the heavy singlet neutrinos via scalar and pseudoscalar couplings, In ref. [8] it was assumed for simplicity a strong hierarchy $\lambda \gg \lambda^{(5)}$ so $\lambda^{(5)}$ was neglected. However, in all the relevant quantities (scatterings, CP asymmetries, light neutrino masses) at leading order the scalar and pseudoscscalar couplings always appear in the combination $\lambda + \lambda^{(5)}$, and thus such an assumption is not necessary. The replacement $\lambda \to \lambda + \lambda^{(5)}$ would suffice to include in the analysis the effects of both type of interactions.

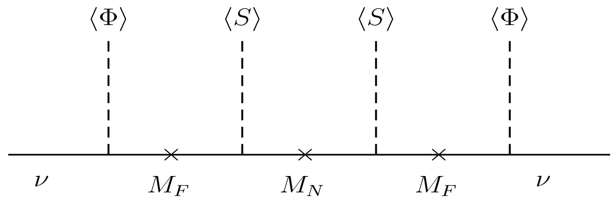

Figure 1. Effective mass operator responsible for neutrino mass generation

2.1. Effective seesaw and light neutrino masses

After $U(1)_X$ and electroweak symmetry breaking the set of Yukawa interactions in (2) generate light neutrino masses through the effective mass operator shown in figure 1. The resulting mass matrix can be written as [8]

$$-\mathcal{M}_{ij} = \left[h^* \frac{\sigma}{M_F} \lambda^T \frac{v^2}{M_N} \lambda \frac{\sigma}{M_F} h^\dagger\right]_{ij} = \left[\tilde{\lambda}^T \frac{v^2}{M_N} \tilde{\lambda}\right]_{ij} . \tag{3}$$

Here we have introduced the seesaw-like couplings

$$\tilde{\lambda}_{\alpha i} = \left(\lambda \frac{\sigma}{M_F} h^\dagger\right)_{\alpha i} . \tag{4}$$

Note that, in contrast to the standard seesaw, the neutrino mass matrix is of fourth order in the *fundamental* Yukawa couplings (h and λ) and due to the factor σ^2/M_F^2 is even more suppressed.

3. Different scenarios for leptogenesis

In this section we discuss the features of each one of the cases we previously mentioned and derive expressions for the CP asymmetries. Henceforth we will use the following notation for the different mass ratios:

$$z_\alpha = \frac{M_{N_\alpha}^2}{M_{N_1}^2}, \qquad \omega_a = \frac{M_{F_a}^2}{M_{F_1}^2}, \qquad r_a = \frac{M_{N_1}}{M_{F_a}}. \tag{5}$$

3.1. The standard leptogenesis case

When the masses of the heavy fields F_a and the $U(1)_X$ symmetry breaking scale are both larger than the Majorana neutrino masses ($M_F\,, \sigma > M_N$) there are no major differences from the standard Fukugita-Yanagida leptogenesis model [4]. After integrating out the F fields one obtains the standard seesaw Lagrangian containing the effective operators $\tilde{\lambda}_{\alpha i} \bar{N}_\alpha l_i \Phi$ with the seesaw couplings $\tilde{\lambda}_{\alpha i}$ given in eq. (4). The right handed neutrino N_1 decays predominantly via 2-body channels as shown in fig. 2. This yields the standard results that for convenience we recall here. The total decay width is $\Gamma_{N_1} = (M_{N_1}/16\pi)\,(\tilde{\lambda}\tilde{\lambda}^\dagger)_{11}$ and the sum of the vertex and self-energy contributions to the CP-asymmetry for N_1 decays into the flavor l_j reads [10]

$$\epsilon_{N_1 \to l_j} = \frac{1}{8\pi (\tilde{\lambda}\tilde{\lambda}^\dagger)_{11}} \sum_{\beta \neq 1} \mathbb{Im} \left\{ \tilde{\lambda}_{\beta j} \tilde{\lambda}_{1j}^* \left[(\tilde{\lambda}\tilde{\lambda}^\dagger)_{\beta 1} \tilde{F}_1(z_\beta) + (\tilde{\lambda}\tilde{\lambda}^\dagger)_{1\beta} \tilde{F}_2(z_\beta) \right] \right\} , \tag{6}$$

where

$$\tilde{F}_1(z) = \frac{\sqrt{z}}{1-z} + \sqrt{z} \left(1 - (1+z) \ln \frac{1+z}{z} \right), \qquad \tilde{F}_2(z) = \frac{1}{1-z}. \tag{7}$$

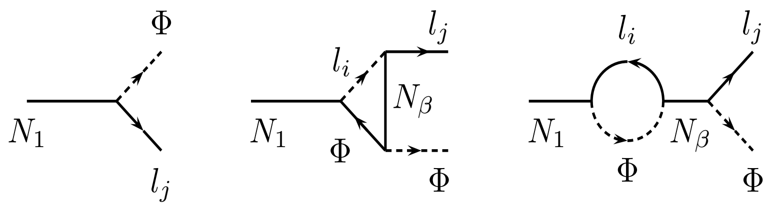

Figure 2. Diagrams responsible for the CP violating asymmetry in the standard case

At leading order in $1/z_\beta$ and after summing over all leptons l_j, eq. (6) yields for the total asymmetry:

$$\epsilon_{N_1} = \frac{3}{16\pi(\tilde{\lambda}\tilde{\lambda}^\dagger)_{11}} \sum_\beta \mathbb{Im}\left\{\frac{1}{\sqrt{z_\beta}}(\tilde{\lambda}\tilde{\lambda}^\dagger)^2_{\beta 1}\right\}. \tag{8}$$

where the sum over the heavy neutrinos has been extended to include also N_1 since for $\beta = 1$ the corresponding combination of couplings is real.

In the hierarchical case $M_{N_1} \ll M_{N_{2,3}}$ the size of the total asymmetry in (8) is bounded by the Davidson-Ibarra limit [11]

$$|\epsilon_{N_1}| \leq \frac{3}{16\pi}\frac{M_{N_1}}{v^2}(m_{\nu_3} - m_{\nu_1}) \lesssim \frac{3}{16\pi}\frac{M_{N_1}}{v^2}\frac{\Delta m^2_{\rm atm}}{2m_{\nu_3}}, \tag{9}$$

where m_{ν_i} (with $m_{\nu_1} < m_{\nu_2} < m_{\nu_3}$) are the light neutrinos mass eigenstates and $\Delta m^2_{\rm atm} \sim 2.5 \times 10^{-3}\,{\rm eV}^2$ is the atmospheric neutrino mass difference [12]. It is now easy to see that (9) implies a lower limit on M_{N_1}. The amount of B asymmetry that can be generated from N_1 dynamics can be written as:

$$\frac{n_B}{s} = -\kappa_s\,\epsilon_{N_1}\,\eta, \tag{10}$$

where $\kappa_s \approx 1.3 \times 10^{-3}$ accounts for the dilution of the asymmetry due to the increase of the Universe entropy from the time the asymmetry is generated with respect to the present time, η (that can range between 0 and 1, with typical values $10^{-1} - 10^{-2}$) is the *efficiency factor* that accounts for the amount of L asymmetry that can survive the washout process. Assuming that ϵ_{N_1} is the main source of the $B - L$ asymmetry [13], eqs. (9) and (10) in addition from the oserved baryon asymmetry eq. (1) yield:

$$M_{N_1} \gtrsim 10^9 \frac{m_{\nu_3}}{\eta\sqrt{\Delta m^2_{\rm atm}}}\,{\rm GeV}. \tag{11}$$

This limit can be somewhat relaxed depending on the specific initial conditions [14] or when flavor effects are included [15, 16, 17, 18] but the main point remains, and that is that the value of M_{N_1} should be well above the electroweak scale.

3.2. Purely flavored leptogenesis case

Differently from standard leptogenesis in the present case, since $M_F > M_{N_1}$, two-body N_1 decays are kinematically forbidden. However, via off-shell exchange of the heavy F_a fields, N_1 can decay to the three body final states $S\Phi l$ and $\bar{S}\bar{\Phi}\bar{l}$. The corresponding Feynman diagram is depicted in figure 3(a). At leading order in $r_a = M_{N_1}/M_{F_a}$, the total decay width reads [8]

$$\Gamma_{N_1} \equiv \sum_j \Gamma(N_1 \to S\Phi l_j + \bar{S}\bar{\Phi}\bar{l}_j) = \frac{M_{N_1}}{192\pi^3}\left(\frac{M_{N_1}}{\sigma}\right)^2 (\tilde{\lambda}\tilde{\lambda}^\dagger)_{11}\,. \tag{12}$$

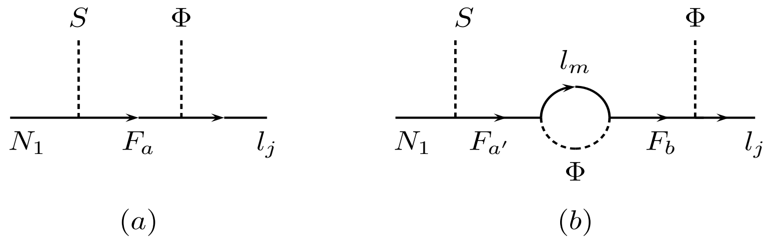

Figure 3. Feynman diagrams responsible for the CPV asymmetry.

As usual, CPV asymmetries in N_1 decays arise from the interference between tree-level and one-loop amplitudes. As was noted in [8], in this model at one-loop there are no contributions from vertex corrections, and the only contribution to the CPV asymmetries comes from the self-energy diagram 3(b). Summing over the leptons and vectorlike fields running in the loop, at leading order in r_a the CPV asymmetry for N_1 decays into leptons of flavor j can be written as

$$\epsilon_{1j} \equiv \epsilon_{N_1 \to \ell_j} = \frac{3}{128\pi} \frac{\sum_m \mathbb{Im}\left[\left(hr^2h^\dagger\right)_{mj} \tilde{\lambda}_{1m}\tilde{\lambda}^*_{1j}\right]}{\left(\tilde{\lambda}\tilde{\lambda}^\dagger\right)_{11}}. \tag{13}$$

Note that since the loop correction does not violate lepton number, the total CPV asymmetry that is obtained by summing over the flavor of the final state leptons vanishes [19], that is $\epsilon_1 \equiv \sum_j \epsilon_{1j} = 0$. This is the condition that defines PFL; namely there is no CPV *and* lepton number violating asymmetry, and the CPV lepton flavor asymmetries are the only seed of the Cosmological lepton and baryon asymmetries.

It is important to note that the effective couplings $\tilde{\lambda}$ defined in eq. (4) are invariant under the reparameterization

$$\lambda \to \lambda \cdot (rU)^{-1}, \quad h^\dagger \to (Ur) \cdot h^\dagger\,, \tag{14}$$

where U is an arbitrary 3×3 non-singular matrix. Clearly the light neutrino mass matrix is invariant under this transformation. Moreover, also the flavor dependent washout processes, that correspond to tree level amplitudes that are determined, to a good approximation, by the effective $\tilde{\lambda}$ couplings, are left essentially unchanged.[2] On the contrary, the flavor CPV asymmetries eq. (13), that are determined by loop amplitudes containing an additional factor of $hr^2h^\dagger$, get rescaled as $hr^2h^\dagger \to h(rUr)^\dagger(rUr)h^\dagger$. Clearly, this rescaling affects in the same way all the lepton flavors (as it should be to guarantee that the PFL conditions $\epsilon_\alpha \equiv \sum \epsilon_{\alpha j} = 0$ are not spoiled), and thus for simplicity we will consider only rescaling by a global scalar factor $r.U = U.r = \kappa\, I$ (with I the 3×3 identity matrix) that, for our purposes, is completely equivalent to the more general transformation (14). Thus, while rescaling the Yukawa couplings through

$$\lambda \to \lambda\, \kappa^{-1}, \quad h^\dagger \to \kappa\, h^\dagger\,, \tag{15}$$

does not affect neither low energy neutrino physics nor the washout processes, the CPV asymmetries get rescaled as:

$$\epsilon_{1j} \to \kappa^2 \epsilon_{1j}\,. \tag{16}$$

By choosing $\kappa > 1$, all the CPV asymmetries get enhanced as κ^2 and, being the Cosmological asymmetries generated through leptogenesis linear in the CPV asymmetries, the final result gets enhanced by the same factor. Therefore, for any given set of couplings, one can always

[2] The approximation is exact in the limit of pointlike F-propagators $(s - M_F^2 + iM_F\Gamma_F) \to M_F^2$.

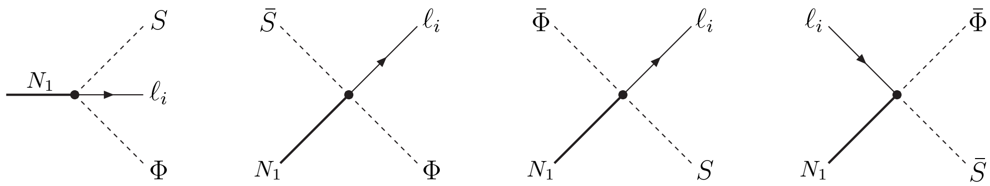

Figure 4. Feynman diagrams for $1 \leftrightarrow 3$ and $2 \leftrightarrow 2$ s, t and u channel processes after integrating out the heavy vectorlike fields F_a.

find an appropriate rescaling such that the correct amount of Cosmological lepton asymmetry is generated. In practice, the rescaling factors κ cannot be arbitrarily large: first, they should respect the condition that all the fundamental Yukawa couplings remain in the perturbative regime; second the size of the h couplings (and thus also of the rescaling parameter κ) is also constrained by experimental limits on lepton flavor violating decays.

3.2.1. Boltzmann Equations In this section we compute the lepton asymmetry by solving the appropriate BE. In general, to consistently derive the evolution equation of the lepton asymmetry all the possible processes at a given order in the couplings have to be included. In the present case $1 \leftrightarrow 3$ decays and inverse decays, and $2 \leftrightarrow 2$ s, t and u channel scatterings all occur at the same order in the couplings and must be included altogether in the BE. The Feynman diagrams for these processes are shown in Figure 4. In addition, the CPV asymmetries of some higher order multiparticle reactions involving the exchange of one off-shell N_1, also contribute to the source term of the asymmetries at the same order in the couplings than the CPV asymmetries of decays and $2 \leftrightarrow 2$ scatterings. More precisely, for a proper derivation of the BE it is essential that the CPV asymmetries of the off-shell $3 \leftrightarrow 3$ and $2 \leftrightarrow 4$ scattering processes are also taken into account [9].

As regards the equation for the evolution of the heavy neutrino density Y_{N_1}, only the diagrams in fig. 4, that are of leading order in the couplings, are important [9]

$$\dot{Y}_{N_1} = -\left(y_{N_1} - 1\right)\gamma_{\rm tot}\,, \tag{17}$$

$$\dot{Y}_{\Delta\ell_i} = \left(y_{N_1} - 1\right)\epsilon_i\gamma_{\rm tot} - \Delta y_i\left[\gamma_i + \left(y_{N_1} - 1\right)\gamma^{N_1\bar{\ell}_i}_{S\Phi}\right]\,, \tag{18}$$

Here we have normalized particle densities to their equilibrium densities $y_a \equiv Y_a/Y_a^{\rm eq}$ where $Y_a = n_a/s$ with n_a the particle number density and s the entropy density. The time derivative is defined as $\dot{Y} = sHz\,dY/dz$ with $z = M_{N_1}/T$ and H is the Hubble parameter. In the last term of the second equation we have used the compact notation for the reaction densities $\gamma^{N_1\bar{\ell}_i}_{S\Phi} = \gamma(N_1\bar{\ell}_i \to S\Phi)$ and in addition we have defined

$$\gamma_i \;=\; \gamma^{N_1}_{S\ell_i\Phi} + \gamma^{N_1\bar{S}}_{\Phi\ell_i} + \gamma^{N_1\bar{\Phi}}_{S\ell_i} + \gamma^{N_1\bar{\ell}_i}_{S\Phi}\,, \tag{19}$$

$$\gamma_{\rm tot} \;=\; \sum_{i=e,\mu,\tau}\gamma_i + \bar{\gamma}_i\,, \tag{20}$$

where in the second equation $\bar{\gamma}_i$ represents the sum of the CP conjugates of the processes summed in γ_i.

Since in this model N_1 decays are of the same order in the couplings than scatterings (that is $\mathcal{O}(\tilde{\lambda}^2)$), the appropriate condition that defines the *strong washout* regime in the case at hand reads:

$$\left.\frac{\gamma_{\rm tot}}{z\,H\,s}\right|_{z\sim 1} > 1 \qquad \text{(strong washout)}, \tag{21}$$

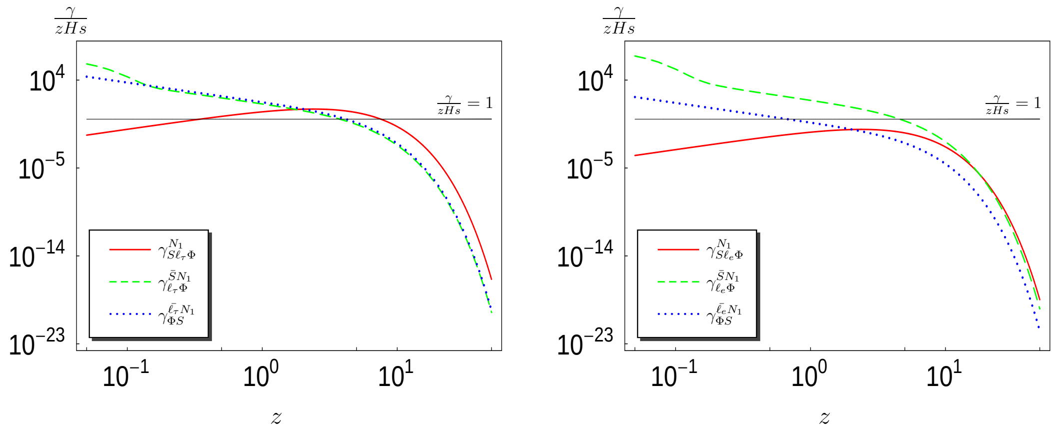

Figure 5. Reaction densities normalized to zHs for $N_1 \to S\ell\Phi$ decays (red solid lines), s-channel $\bar{S}N_1 \leftrightarrow \ell\Phi$ scatterings (green dashed lines), and t, u-channel scatterings in the point-like approximation (blue dotted lines). Left panel: τ flavor. Right panel: electron flavor.

and conversely $\gamma_{\rm tot}/(z\,H\,s)|_{z\sim 1} < 1$ defines the *weak washout* regime. Note that this is different from standard leptogenesis, where at $z \sim 1$ two body decays generally dominate over scatterings, and e.g. the condition for the strong washout regime can be approximated as $\gamma_{\rm tot}/(z\,H\,s)|_{z\sim 1} \sim \Gamma_{N_1}/H|_{z\sim 1} > 1$.

3.2.2. Results In this section we discuss a typical example of successful leptogenesis at the scale of a few TeV. The example presented is a general one. No particular choice of the parameters has been performed, except for the fact that the low energy neutrino data are reproduced within errors, and that the choice yields an interesting washout dynamics well suited to illustrate how PFL works. The numerical value of the final lepton asymmetry ($Y_{\Delta L} \sim -7.2 \times 10^{-10}$) is about a factor of 3 *larger* than what is indicated by measurements of the Cosmic baryon asymmetry. This is however irrelevant since, as was already discussed, it would be sufficient a minor rescaling of the couplings (or a slight change in the CPV phases) to obtain the precise experimental result. In the numerical analysis we have neglected the dynamics of the heavier singlet neutrinos since the N_α masses are sufficiently hierarchical to ensure that $N_{2,3}$ related washouts do not interfere with N_1 dynamics. Moreover, in the (strong washout) fully flavored regime (that is effective as long as $T < 10^9\,$GeV) the $N_{2,3}$ CPV asymmetries do not contribute to the final lepton number asymmetry [13].

In figure 5 we show the behavior of the various reaction densities for decays and scatterings, normalized to sHz, as a function of z. The results correspond to a mass of the lightest singlet neutrino fixed to $M_{N_1} = 2.5\,$TeV, the heavier neutrino masses are $M_{N_2} = 10\,$TeV and $M_{N_3} = 15\,$TeV, and the relevant mass ratios $r_a = M_{N_1}/M_{F_a}$ for the messenger fields are $r_{1,2,3} = 0.1, 0.01, 0.001$ (the effects of the lightest F resonances can be seen in the s-channel rates in both panels in fig. 5). The fundamental Yukawa couplings h and λ are chosen to satisfy the requirement that the seesaw formula eq. (3) reproduces within $2\,\sigma$ the low energy data on the neutrino mass squared differences and mixing angles [12]. Typically, when this requirement is fulfilled, one also ends up with a dynamics for all the lepton flavors in the strong washout regime. This is shown in the left panel in fig. 6 where we present the total rates for the three flavors.

The left panel in fig. 5 refers to the decay and scattering rates involving the τ-flavor that, in our example, is the flavor more strongly coupled to N_1, and that thus suffers the strongest

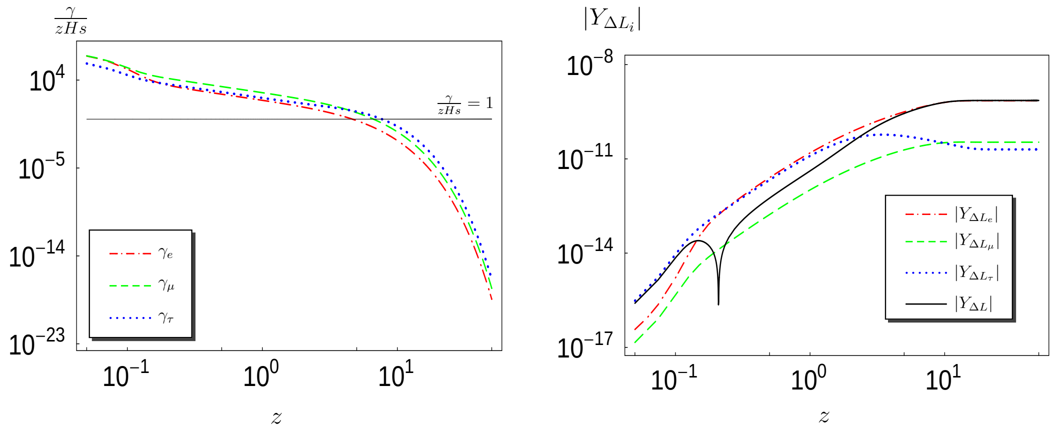

Figure 6. Left panel: the total washout rates for each lepton flavor normalized to zHs as a function of z. Right panel: the evolution of the absolute value of the flavored density asymmetries and of the lepton number asymmetry (black solid line). The flavor CPV asymmetries are $\epsilon_{1e} = -4.7 \times 10^{-4}$, $\epsilon_{1\mu} = -1.9 \times 10^{-4}$ and $\epsilon_{1\tau} = 6.6 \times 10^{-4}$. The final values of the asymmetry densities (at $z \gg 1$) are $Y_{\Delta_{\ell_e}} = -7.1 \times 10^{-10}$, $Y_{\Delta_{\ell_\mu}} = -0.3 \times 10^{-10}$, $Y_{\Delta_{\ell_\tau}} = 0.2 \times 10^{-10}$.

washout. It is worth noticing that, due to the fact that in this model scatterings are not suppressed by additional coupling constants with respect to the decays, the decay rate starts dominating the washouts only at $z \gtrsim 1$. The right panel in fig. 5 depicts the reaction rates for the electron flavor, that is the more weakly coupled, and for which the strong washout condition eq. (21) is essentially ensured by sizeable s-channel scatterings. Scatterings and decay rates for the μ-flavor are not shown, but they are in between the ones of the previous two flavors.

The total reaction densities that determine the washout rates for the different flavors are shown in the first panel in figure 6. The evolution of these rates with z should be confronted with the evolution of (the absolute value of) the asymmetry densities for each flavor, depicted in the second panel on the right. Since, as already stressed several times, PFL is defined by the condition that the sum of the flavor CPV asymmetry vanishes ($\sum_j \epsilon_{1j} = 0$), it is the hierarchy between these washout rates that in the end is the responsible for generating a net lepton number asymmetry. In the case at hand, the absolute values of the flavor CPV asymmetries satisfy the condition $|\epsilon_\mu| < |\epsilon_e| < |\epsilon_\tau|$, as can be inferred directly by the fact that at $z < 0.1$, when the effects of the washouts are still negligible, the asymmetry densities satisfy this hierarchy. Moreover, since $\epsilon_{\mu,e} < 0$ while $\epsilon_\tau > 0$, initially the total lepton number asymmetry, that is dominated by $Y_{\Delta_{\ell_\tau}}$, is positive. As washout effects become important, the τ-related reactions (blue dotted line in the left panel) start erasing $Y_{\Delta_{\ell_\tau}}$ more efficiently than what happens for the other two flavors, and thus the initial positive asymmetry is driven towards zero, and eventually changes sign around $z = 0.2$. This change of sign corresponds to the steep valley in the absolute value $|Y_{\Delta L}|$ that is drawn in the figure with a black solid line. Note that when all flavors are in the strong washout regime, as in the present case, the condition for the occurrence of this 'sign inversions' is simply given by $\max_{j\in e,\mu}\left(|\epsilon_j|/|\tilde{\lambda}_{1j}|^2\right) \gtrsim \epsilon_\tau/|\tilde{\lambda}_{1\tau}|^2$. From this point onwards, the asymmetry remains negative, and since the electron flavor is the one that suffers the weakest washout, $Y_{\Delta_{\ell_e}}$ ends up dominating all the other density asymmetries. In fact, as can be seen from the right panel in fig. 6, it is $Y_{\Delta_{\ell_e}}$ that determines to a large extent the final value of the lepton asymmetry $Y_{\Delta L} = -7.2 \times 10^{-10}$.

A few comments are in order regarding the role played by the F_a fields. Even if $M_{N_1} \ll M_{F_a}$, at large temperatures $z \gg 1$ the tail of the thermal distributions of the N_1, S and Φ particles allows the on-shell production of the lightest F states. A possible asymmetry generated in the decays of the F fields can be ignored for two reasons: first because due to the rather large h and λ couplings F decays occur to a good approximation in thermal equilibrium, ensuring that no sizeable asymmetry can be generated, and second because the strong washout dynamics that characterizes N_1 leptogenesis at lower temperatures is in any case insensitive to changes in the initial conditions.

In conclusion, it is clear from the results of this section that the model encounters no difficulties to allow for the possibility of generating the Cosmic baryon asymmetry at a scale of a few TeVs. Moreover, our analysis provides a concrete example of PFL, and shows that the condition $\epsilon_1 \neq 0$ is by no means required for successful leptogenesis.

4. Conclusions

Variations of the standard leptogenesis picture can arise from the presence of an additional energy scale different from that of lepton number violation. Quite generically the resulting scenarios are expected to yield qualitative and quantitative changes on leptogenesis. Here we have considered what we regard as the simplest possibility namely, the presence of an abelian flavor symmetry $U(1)_X$. We have described two possible scenarios within this framework and have explored their implications for leptogenesis.

We have found that as long as the abelian flavor symmetry energy scales remain above the lepton number violating scale neither qualitative nor quantitative differences with the standard leptogenesis model arise. Conversely if the $U(1)_X$ is unbroken during the leptogenesis era and the messengers fields F_a are too heavy to be produced on-shell in N_1 decays *purely flavored leptogenesis* at the TeV scale results. By solving the corresponding BE we have shown that within this scenario the non-vanishing of the CPV lepton flavor asymmetries in addition to the lepton and flavor violating washout processes occuring in the plasma provide the necessary ingredients to generate the Cosmic baryon asymmetry. Accordingly, if below the leptogenesis scale new energy scales are present -as might be expected- the interplay between these scales could have a quite interesting impact on leptogenesis.

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