

Frascati Physics Series Vol. XLVIII (2009), pp. 13-18
YOUNG RESEARCHERS WORKSHOP: “Physics Challenges in the LHC Era”
Frascati, May 11 and 14, 2009

PURELY FLAVORED LEPTOGENESIS AT THE TeV SCALE

Luis Alfredo Muñoz

Instituto de Física, Universidad de Antioquia, A.A. 1226, Medellín, Colombia

Abstract

I study variations of the standard leptogenesis scenario that can arise if an additional mass scale related to the breaking of some new symmetry is present below the mass M_{N_1} of the lightest right-handed Majorana neutrino. I present a particular realization of this scheme that allows for leptogenesis at the TeV scale. In this realization the baryon asymmetry is exclusively due to flavor effects.

1 Introduction

From observations of light element abundances and of the Cosmic microwave background radiation ¹⁾ the Cosmic baryon asymmetry, $Y_B = \frac{n_B - n_{\bar{B}}}{s} = (8.75 \pm 0.23) \times 10^{-10}$, (where s is the entropy density) 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 throughout discussion). Leptogenesis ⁴⁾ is a scenario in which an initial lepton asymmetry, generated in the out-of-equilibrium decays of heavy 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 has been achieved (see ref. ⁷⁾).

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 ($N_\alpha = N_{\alpha R} + N_{\alpha R}^c$) 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 $U(1)_X$ gauge 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

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

We use Greek indices $\alpha, \beta \dots = 1, 2, 3$ to label the heavy Majorana neutrinos, Latin indices $a, b \dots = 1, 2, 3$ for the vectorlike messengers, and i, j, k, \dots for the lepton flavors e, μ, τ . Following reference ⁸⁾ we chose the simple $U(1)_X$ charge assignments $X(\ell_{Li}, F_{La}, F_{Ra}) = +1$, $X(S) = -1$ 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. As discussed in ⁸⁾, depending on the hierarchy between the relevant scales of the model

$(M_{N_1}, M_{F_a}, \sigma)$, quite different *scenarios* for leptogenesis can arise: *i*) For $M_F, \sigma \gg M_N$, we recover the Standard Leptogenesis (SL) case that will not discuss. *ii*) For $\sigma < M_{N_1} < M_{F_a}$ we obtain the Purely Flavored Leptogenesis (PFL) case, that corresponds to the situation, 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 (for other possibilities see ref. ⁸). After $U(1)_X$ and electroweak symmetry breaking the set of Yukawa interactions in (1) generates light neutrino masses through the effective mass operator. The resulting mass matrix can be written as ⁸)

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

Here we have introduced the seesaw-like couplings $\tilde{\lambda}_{\alpha i} = \left(\lambda \frac{\sigma}{M_F} h^\dagger \right)_{\alpha i}$. 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 Purely flavored leptogenesis

In the case when $\sigma < M_{N_1} < M_F$, 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 CP asymmetry is obtained from the interference between the tree level and loop diagrams; and reads ⁸):

$$\epsilon_{N_1 \rightarrow \ell_j} \equiv \epsilon_j = \frac{3}{128\pi} \frac{\sum_i \Im m \left[(hr^2 h^\dagger)_{ij} \tilde{\lambda}_{1i} \tilde{\lambda}_{1j}^* \right]}{\left(\tilde{\lambda} \tilde{\lambda}^\dagger \right)_{11}}, \quad (3)$$

where $r = M_{N_1}/M_F$. The CP asymmetries in (3) have the following properties: *i*) $\epsilon_j \neq 0$. *ii*) The total CP asymmetry $\epsilon_{N_1} = \sum_j \epsilon_j = 0$. This is because $\Im m[\tilde{\lambda} h r^2 h^\dagger \tilde{\lambda}^\dagger]_{11} = 0$, and is related to the fact that the loop does not involve lepton number violation. *iii*) Rescaling the couplings h and λ by a parameter $\kappa > 1$ according to: $h \rightarrow \kappa h$, $\lambda \rightarrow \kappa^{-1} \lambda$; enhances the CP asymmetries as $\epsilon_i \rightarrow \kappa^2 \epsilon_i$ ⁹). In the Boltzmann Equations (BE), the particles densities are written in terms of the entropy density s , i.e $Y_a = n_a/s$ where n_a is the number density for the particle a . We rescale the densities Y_a by the equilibrium density

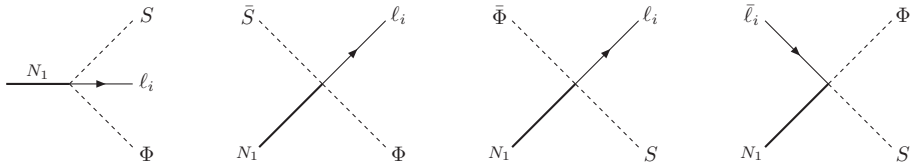


Figure 1: Relevant diagrams for leptogenesis in our PFL model.

Y_a^{eq} of the corresponding particle, defining $y_a \equiv Y_a/Y_a^{eq}$, while the asymmetries of the rescaled densities are denoted by $\Delta y_a \equiv y_a - y_{\bar{a}}$. In the BE, the time derivative is defined as $\dot{Y}_a = z H s \frac{dY_a}{dz}$ where $z = M_{N_1}/T$ and $H(z)$ is the Hubble rate at temperature T . We denote the thermally averaged rate for an initial state A to go the final state B as $\gamma_B^A = \gamma(A \rightarrow B)$. The processes in the BE are the decay of N_1 and the scatterings illustrated in figure 1; they are all of the same order $\mathcal{O}((\lambda^\dagger h)^2)$. The BE for the evolution of the N_1 abundance and of the lepton density asymmetry $Y_{\Delta L_i}$ are (see ⁹) for details):

$$\begin{aligned} \dot{Y}_{N_1} &= -(y_{N_1} - 1) \gamma \\ \dot{Y}_{\Delta L_i} &= (y_{N_1} - 1) \epsilon_i \gamma - \Delta y_i \left(\gamma_i + (y_{N_1} - 1) \gamma_{S\Phi}^{N_1 \bar{\ell}_i} \right), \end{aligned} \quad (4)$$

where $\gamma_i = \gamma_{S\Phi \ell_i}^{N_1} + \gamma_{\Phi \ell_i}^{\bar{S} N_1} + \gamma_{S \ell_i}^{\bar{\Phi} N_1} + \gamma_{S\Phi}^{\bar{\ell}_i N_1}$ is the sum of the processes depicted in fig. 1 and $\gamma = \sum_i \gamma_i + \bar{\gamma}_i$. For PFL the strong washout condition, corresponds to: $\frac{\gamma}{z H s} \Big|_{z \sim 1} > 1$ where the normalization factor $z H s$ has been chosen to obtain an adimensional ratio. For the BE's solution we have chosen the couplings h and λ such that they reproduce the low energy neutrino parameters within 2σ and also satisfy the strong washout condition. The densities rates for e and τ are shown in figure 2, their expressions can be found in ⁹). It is important to remark that the value of the lowest Majorana mass has been chosen at the TeV scale: $M_{N_1} = 2.5$ TeV. Others parameters are $M_{N_2} = 10$ TeV, $M_{N_3} = 15$ TeV and $r_a = [.1, .01, .001]$. The total reaction densities that determine the washout rates for the different flavors are shown in the first panel in figure 3. Since PFL is defined by the condition that the sum of the flavor CP 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. The total lepton asymmetry obtained is above the experimental data, but the couplings can be rescaled by a suitable value of κ to get the right value.

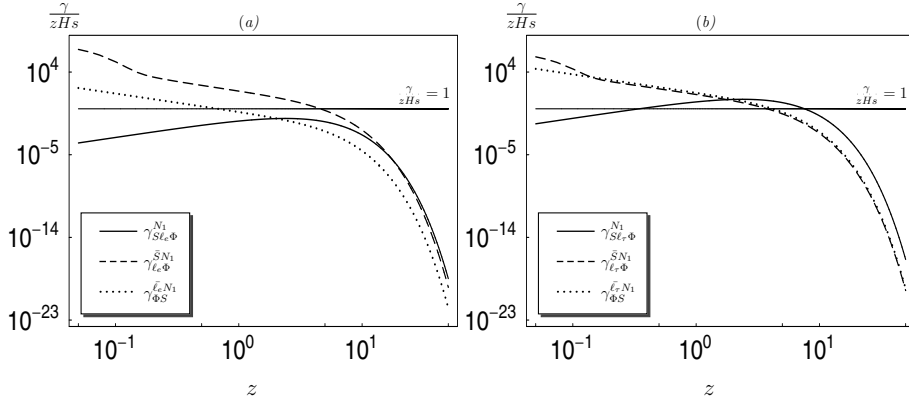


Figure 2: (a) The e density rates. Around $z \sim 1$ s -channel scatterings is the dominant one. (b) The τ density rates. Around $z \sim 1$ both decay and scatterings are faster than $z H s$.

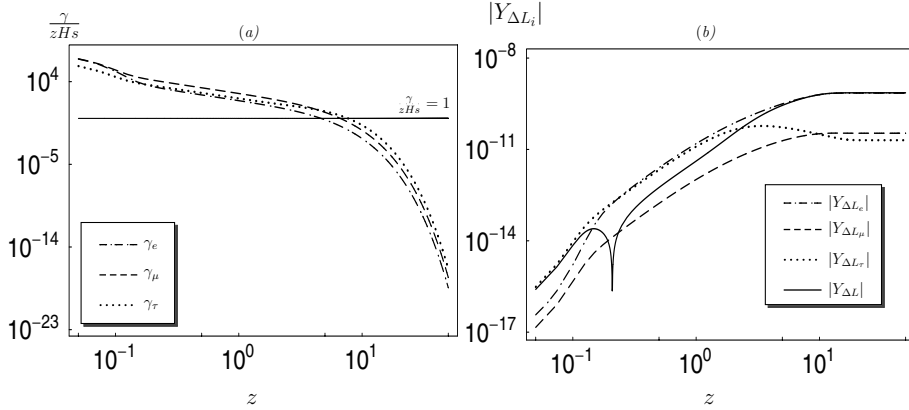


Figure 3: (a) Washout rates: with our choice of parameters the washout for the τ is stronger (weaker) than that of e and μ , for $z \gg 1$ ($z \ll 1$). (b) Evolution of the flavored and of the total lepton asymmetry. The final value is $Y_{\Delta L} = -7.1 \times 10^{-10}$, that is approximately three times larger than the experimental value, and is completely dominated by $Y_{\Delta L_e}$.

Conclusions

Variations of the standard leptogenesis scenario can arise from the presence of an additional energy scale different from that of lepton number violation.

Quite generically, the resulting scenarios can be expected to be qualitatively and quantitatively different from SL. Here we have considered what we regard as the simplest possibility namely, the presence of an Abelian flavor symmetry $U(1)_X$. The model allows 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.

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