A model for leptogenesis at the TeV scale

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Abstract

We consider the mechanism of thermal leptogenesis at the TeV scale in the context of an extension of the Standard Model with 4 generations and the inclusion of four right-handed Majorana neutrinos. We obtain a value for the baryon asymmetry of the Universe in accordance with observations by solving the full set of coupled Boltzmann equations for this model.

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1 Introduction

The attractive mechanism of thermal leptogenesis has been thoroughly investigated in recent years. Most of the focus has been on the extended model of the Standard Model (SM) with 3 right-handed (RH) neutrinos \[1\]. This simple model which can explain the baryon asymmetry of the Universe (BAU) and provide appropriate values for neutrino masses and mixing angles is based on the out-of-equilibrium decay of a RH neutrino, in a CP violating way, such that an asymmetry in the leptonic decay products is produced. The leptonic asymmetry is then converted into a baryonic asymmetry due to \((B+L)\)-violating sphaleron interactions which are in equilibrium above the electroweak breaking scale.

From the experimental point of view the recent best fit values of the ratio of the baryon to photon density \(\eta = \frac{n_B}{n_\gamma}\)

\[
\eta = 6.5^{+0.4}_{-0.3} \times 10^{-10}.
\]

is given by WMAP [2].

The typical scale at which the standard scenario described above occurs corresponds to \(10^{10} - 10^{15}\) GeV, and the induced neutrino masses arise via the see-saw mechanism \[3, 4, 5\] with a common interaction being at the origin of the asymmetry and mass generation.

Recently, detailed calculations have been performed solving the corresponding Boltzmann equations of the system taking into account the different interactions which contribute both to the production and potential washout of the leptonic asymmetry. The implications of these analyses have provided strong constraints on the values of the light left-handed neutrino masses \(\sum m_i \sim 0.2\) eV \[6, 7, 8\] and a lower bound on the mass of the lightest RH neutrino \(M \gtrsim 10^8\) GeV \[6, 9, 10\]. The more recent work of ref. \[11\] has included the effect of scattering with gauge bosons which had been neglected before and in ref. \[12\] a comprehensive calculation including effects from renormalization group (RG) running, finite temperature background effects, other possible decays producing a leptonic asymmetry (decay of a scalar particle), a correction to the appropriate subtraction that must be performed for scattering via N-exchange was presented. A thorough recalculation of this last issue was done in ref. \[13\], coinciding with ref. \[12\] that the washout contribution from this type of diagram had been overestimated.

The less standard scenario of leptogenesis occurring at the TeV scale has been investigated for example in references \[14, 15, 16, 17, 18, 19, 20, 21\]. In particular, in ref. \[14\] a TeV scale model which is a simple extension of the Fukugita-Yanagida (FY) model including one more generation to the SM and four gauge singlets was discussed and estimates of the produced CP-asymmetry and final baryon asymmetry were obtained. Both supersymmetric and non-supersymmetric versions of the model were analyzed.

In this paper we focus on this model of ref. \[14\] and solve the corresponding Boltzmann equations (BE) in order to reliably obtain the final values of the asymmetry and determine the region of parameter space which is available for the model.

We dedicate section 2 to a careful discussion of the way we set up our calculations of the Boltzmann equations making emphasis on a few subtle points. In section 3, we recall the main features of the model and establish our notation. Section 4 presents the Boltzmann equations valid for this model. Section 5 is devoted to our results and conclusions.
2 Preliminary Considerations

In reference \[22\] it was pointed out that the relevant object that should be studied in thermal leptogenesis is the density matrix $\rho$. In the FY model $\rho$ is a $(3 \times 3)$ matrix with all entries different from zero for temperatures greater than $10^{12}\text{GeV}$. However, as the temperature decreases and each of the different interactions defined by a specific charged Yukawa coupling enters equilibrium the corresponding off-diagonal elements of the density matrix go to zero. In our model at the TeV scale all of the charged Yukawa couplings are in equilibrium and thus $\rho$ will be a diagonal matrix. The matrix $\rho$ is normalized such that,

$$Tr\rho = \sum_\alpha Y_{L\alpha} \equiv \sum_\alpha (Y_{\ell\alpha} - Y_{\bar{\ell}\alpha}) .$$  \hspace{1cm} (2)

where the sum over $\alpha$ indicates a sum over flavour.

Another important issue is the relationship between the produced lepton asymmetry and the induced $B-L$ asymmetry which is conserved by sphaleron interactions. Again a careful consideration of the temperatures and which interactions are in equilibrium changes the equation relating these asymmetries, see reference \[22\]. In general the procedure that is employed is through the establishment of equations for the chemical potentials reflecting which interactions are in equilibrium. Thus the chemical equilibrium equations allow us to write the asymmetries in number densities in terms of some specific asymmetry, say $Y_{L\alpha}$. For the Fukugita-Yanagida model then generically there would be a $(3 \times 3)$ matrix $A_{\alpha\beta}$ of the form \[22\]

$$Y_{B/3-L\alpha} = A_{\alpha\beta} Y_{L\beta} ,$$ \hspace{1cm} (3)

and the numerical values for the elements of $A$ are determined from the equations for the chemical potentials. If all interactions related to charged Yukawa couplings are in equilibrium, as well as sphaleron interactions, the equation for the case of $N$ generations simplifies to,

$$\frac{Y_B}{N} - Y_{L\alpha} = \left( -\frac{4}{3N} + \frac{8}{34N+2} \right) \sum_\beta Y_{L\beta} - 3Y_{L\alpha},$$ \hspace{1cm} (4)

or adding over all generations

$$Y_{B-L} = - \left( \frac{22N + 13}{6N + 3} \right) \sum_\alpha Y_{L\alpha} .$$ \hspace{1cm} (5)

Some comments are in order. In the standard scenario of the FY model with three RH neutrinos, where the production of the asymmetry occurs at some high scale for which not all charged Yukawa couplings are in equilibrium, one should use eq. \[3\] to relate the asymmetries. The point is that usually what is solved in the system of coupled Boltzmann equations, under certain assumptions, is an equation for $Y_L = \sum_\alpha Y_{L\alpha}$, see below for a discussion on this point, and then the asymmetry in $B-L$ is calculated using eq. \[5\]. From our previous discussion we see that it is not quite correct to do this unless some approximations are done concerning the different $Y_{L\alpha}$.
For our model at the TeV scale, all charged Yukawa couplings are in equilibrium and it would be correct to use eqns. (4) and (5). Another feature of our model is that the Yukawa coupling of the heavy LH neutrino is of the order of 1 to ensure that the masses of the fourth generation leptons are heavy enough, and also to give an enhancement to the value of the CP-asymmetry produced in the model. However, then the scattering interaction which is associated to the Yukawa coupling of the heavy LH neutrino will also be in equilibrium in our model. This changes the chemical equilibrium equations such that the relationship between $B - L$ ($B$) and $L$ is modified to be:

$$Y_{B-L} = -\left(\frac{13}{3} + \frac{8}{3}(13-3N)\right) \sum_{\alpha=e,\mu,\tau} Y_{L\alpha} = -\left(\frac{22N + 13}{6N - 1}\right) \sum_{\alpha=e,\mu,\tau} Y_{L\alpha}.$$  

(6)

Thus, as was estimated in reference [14] there is no net asymmetry produced in the heavy left-handed lepton. Consequently, we can put $Y_{L\sigma}$ to zero and make the usual assumptions about the other $Y_{L\alpha}$ when constructing the Boltzmann equation for $Y_L$ provided we use eq. (6) to relate the produced leptonic asymmetry to the final net $B - L$ asymmetry. Below, and from now on, $Y_L$ denotes the sum over the asymmetries of only the three light leptonic flavours.

In reference [12] the effect of renormalization group (RG) running for masses and couplings was included. In particular, one of the important effects is that the diagrams with $\Delta L = 1$ proportional to gauge couplings can become sizeable compared to those proportional to the top Yukawa coupling when considering values for these couplings at the high scale at which the lepton asymmetry is being produced. In our model at the TeV scale, we have checked that this is a very small effect and we will not include RG effects in our analysis. We also have not included the $\Delta L = 1$ proportional to gauge couplings as we are not in the resonant leptogenesis case of ref. [11].

Also in ref. [12] thermal background effects have been included into the calculation, it was pointed out that numerically these effects are particularly relevant when correcting the propagator of the Higgs scalar field. We include only this thermal correction into our calculation.

Another important point that should be mentioned, which is crucial for our model, is the fact that the actual decay temperature at which the lepton asymmetry is produced is just above the electroweak scale; thus one must be careful and check that for our choice of parameters the sphaleron interactions are still in equilibrium. Although, in the SM there is no first order phase transition, just a crossover, and this is not modified for our model with four generations, sphaleron interactions are switched off for temperatures below $\sim 100$ GeV [24]. A careful calculation of the chemical equilibrium equations in the region close to the electroweak phase transition was performed in [25].

3 The Model

The relevant part of the Lagrangian of the model we consider is given by,

$$L = L_{SM} + \bar{\psi}_R i \partial \psi_R - \frac{M_{N_i}}{2} (\bar{\psi}_R \psi_R + h.c.) - (\lambda^{\nu}_{\alpha} \bar{L}_i \psi_{R_i} \phi + h.c.),$$  

(7)

1In this model, we denote by $\sigma$ the flavour of the fourth leptonic generation.
where $\psi_{R_i}$ are two-component spinors describing the right-handed neutrinos and we define a Majorana 4-component spinor, $N_i = \psi_{R_i} + \psi^c_{R_i}$. Our index $i$ runs from 1 to 4, and $\alpha = e, \mu, \tau, \sigma$. The $\sigma$ component of $L_\alpha$ corresponds to a left-handed lepton doublet which must satisfy the LEP constraints from the $Z$-width on a fourth left-handed neutrino [26]. The $\lambda_{\nu_i}^\alpha$ are Yukawa couplings and the field $\phi$ is the SM Higgs boson doublet whose vacuum expectation value is denoted by $v$.

We work in the basis in which the mass matrix for the right-handed neutrinos $M$ is diagonal and real,

$$M = \text{diag}(M_1, M_2, M_3, M_4)$$

and define $m_D = \lambda^\nu v$.

To leading order $^2$, the induced see-saw neutrino mass matrix for the left-handed neutrinos is given by,

$$m_\nu = \lambda^{\nu\dagger}M^{-1}\lambda^\nu v^2.$$ 

which is diagonalized by the well known PMNS matrix.

We consider the out-of-equilibrium decay of the lightest of the gauge singlets $N_j$, which we take to be $N_1$. The decay rate at tree-level is given by

$$\Gamma_{N_j} = \frac{(\lambda^{\dagger}\lambda)_{jj}}{8\pi} M_j = \sum_{\alpha=1}^4 \frac{(\lambda_{\alpha j}^\dagger\lambda_{\alpha j})}{8\pi} M_j = \frac{\tilde{m}_j^{(4G)} M_1 M_j}{8\pi v^2},$$

where $\lambda \equiv \lambda^\nu$ and we have defined $^3$

$$\tilde{m}_j^{(4G)} = \sum_{\alpha=1}^4 \frac{(\lambda_{\alpha j}^\dagger\lambda_{\alpha j})}{M_1} v^2.$$ 

To ensure an out-of-equilibrium decay of $N_1$ it is necessary that $\frac{\Gamma_{N_1}}{H(T=M_1)} \ll 1$, where $H$ is the Hubble expansion rate at $T = M_1$. This condition will impose an upper bound on the usual effective mass parameter $\tilde{m}_1$ which is defined in for example [27].

The CP asymmetry calculated from the interference of the tree diagrams with the one-loop diagrams (self-energy and vertex corrections) is [28]

$$\epsilon_i = \frac{1}{(8\pi)} \frac{1}{|\lambda_{\alpha_i}^\nu\lambda_{\alpha_i}^\nu|_{ii}} \sum_j \text{Im}[\lambda^{\nu\dagger} \lambda_{\nu i}]^2 \left[ f \left( \frac{M_j^2}{M_i^2} \right) + g \left( \frac{M_j^2}{M_i^2} \right) \right],$$

$^2$Given the fact that in our model, at least one of the RH neutrino masses take values $\sim \text{TeV}$, one has to go to next order in the general see-saw formula. However, these corrections are small and not relevant for our purpose.

$^3$To simplify our notation, we write the sums over $\alpha = 1$ to 4 denoting the sums over all flavours $e, \mu, \tau, \sigma$, and sums over $\alpha = 1$ to 3 to denote sums over the three lightest flavours. For instance, $\tilde{m}_j^{(3G)} = \sum_{\alpha=1}^3 \frac{(\lambda_{\alpha j}^\dagger\lambda_{\alpha j})}{M_1} v^2$ is the usual effective mass parameter $\tilde{m}_1$ defined in the literature for $j = 1$. 

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where

\[ f(x) = \sqrt{x} \left[ 1 - (1 + x) \ln \frac{1 + x}{x} \right], \]

\[ g(x) = \frac{\sqrt{x}}{1 - x}. \] (13)

In reference [14] it was shown that the upper bound of the CP-asymmetry produced in the decay of the lightest right-handed neutrino \( N_1 \) is

\[ |\epsilon_1| \lesssim \frac{3}{8\pi} \frac{M_1 m_4}{v^2}, \] (14)

where \( m_4 \) denotes the largest eigenvalue of the left-handed neutrino mass matrix \( m_\nu \). Due to experimental constraints we require that \( m_4 > 45 \text{ GeV} \), which implies that the bound on \( \epsilon_1 \) is irrelevant. In general there will be regions of parameter space in which the produced CP asymmetry can be very large, thus the final allowed region of parameter space could include areas in which the washout processes can be very large as well. In reference [14] possible textures for the \( \lambda' \) matrix were given to illustrate how to obtain masses consistent with low energy data and the appropriate amount of CP asymmetry from the decay of \( N_1 \).

4 Boltzmann Equations

We will now write the Boltzmann equations for the RH neutrino abundance and the lepton densities. With all the above arguments in consideration, the main processes we consider in the thermal bath of the early universe are decays, inverse decays of the RH neutrinos, and the lepton number violation \( \Delta L = 1 \) and \( \Delta L = 2 \), Higgs and RH neutrinos exchange scattering processes, respectively.

In our analysis we stick to the case where the asymmetry is due only to the decay of the lightest RH neutrino \( N_1 \). The first BE, which corresponds to the evolution of the abundance of the lightest RH neutrino \( Y_{N_1} \) involving the decays, inverse decays and \( \Delta L = 1 \) processes is given by

\[ \frac{dY_{N_1}}{dz} = -\frac{z}{sH(M_1)} \left( \frac{Y_{N_1}}{\overline{Y}_{N_1}} - 1 \right) \left( \gamma_{D_1} + \gamma_{S_1} \right), \] (15)

where \( z = \frac{M_1}{T} \), \( s \) is the entropy density and \( \gamma_{D_j}, \gamma_{S_j} \) are the interaction rates for the decay and \( \Delta L = 1 \) scattering contributions, respectively.

The BE for the lepton asymmetry is given by

\[ \frac{dY_L}{dz} = -\frac{z}{sH(M_1)} \left( \epsilon_1 \gamma_{D_1} \left( \frac{Y_{N_1}}{\overline{Y}_{N_1}} - 1 \right) + \gamma_{W} \frac{Y_L}{\overline{Y}_{L}} \right), \] (16)

where \( \epsilon_1 \) is the CP violation parameter given by eq. [12] and \( \gamma_{W} \), is a function of \( \gamma_{D_j} \) and \( \gamma_{S_j} \) and \( \Delta L = 2 \) interaction rate processes, called the washout factor which is responsible for the

\(^4\text{We have checked that the asymmetry produced by the decay of the second lightest RH neutrino is very tiny in our model compared to the one produced by the decay of } N_1.\)
damping of the produced asymmetry. In eqs. (15) and (16), \( Y_i^{eq} \) is the equilibrium number density of a particle \( i \), which has a mass \( m_i \), given by

\[
Y_i^{eq}(z) = \frac{45}{4\pi^4} \frac{g_i}{g_*} \left( \frac{m_i}{M_1} \right)^2 z^2 K_2 \left( \frac{m_i z}{M_1} \right),
\]

where \( g_i \) is the internal degree of freedom of the particle \( (g_{N_i} = 2, \ g_\ell = 4) \) and \( g_* \approx 108 \). Explicit expressions of the interaction rates \( \gamma_{D_j} \), \( \gamma_{S_j} \) and \( \gamma_W \) are given below.

The reaction density \( \gamma_{D_j} \) is related to the tree level total decay rate of the RH neutrino \( N_j \) of eq. (10) by

\[
\gamma_{D_j} = \frac{\eta_{eq}^{N_j}}{K_2(z)} \Gamma_{N_j},
\]

where \( K_n(z) \) are the \( K \) type Bessel functions.

We define the reaction density of any process \( a + b \rightarrow c + d \) by

\[
\gamma^{(i)} = \frac{M_1^4}{64\pi^4} \frac{1}{z} \int_0^{\infty} dx \, \frac{\sigma^{(i)}(x) \sqrt{x}}{K_1(\sqrt{x}z)},
\]

where \( \sigma^{(j)}(x) \) are the reduced cross sections for the different processes which contribute to the Boltzmann equations. For the \( \Delta L = 1 \) Higgs boson exchange processes, we have

\[
\sigma_{ij}^{(1)} = 3\alpha_u \sum_{\alpha=1}^{4} \left( \lambda_{\alpha j}^* \lambda_{\alpha j} \right) \left( \frac{x - a_j}{x} \right)^2 = 3\alpha_u \tilde{m}_j^{(4G)} M_1 \left( \frac{x - a_j}{x} \right)^2,
\]

\[
\sigma_{ij}^{(2)} = 3\alpha_u \sum_{\alpha=1}^{4} \left( \lambda_{\alpha j}^* \lambda_{\alpha j} \right) \left( \frac{x - a_j}{x} \right) \left[ \frac{x - 2a_j + 2a_h}{x - a_j + a_h} + \frac{a_j - 2a_h}{x - a_j} \ln \left( \frac{x - a_j + a_h}{a_h} \right) \right]
\]

\[
= 3\alpha_u \tilde{m}_j^{(4G)} M_1 \left( \frac{x - a_j}{x} \right) \left[ \frac{x - 2a_j + 2a_h}{x - a_j + a_h} + \frac{a_j - 2a_h}{x - a_j} \ln \left( \frac{x - a_j + a_h}{a_h} \right) \right],
\]

where

\[
\alpha_u = \frac{Tr(\lambda_u^u \lambda_u)}{4\pi} \simeq \frac{m_u^2}{4\pi v^2}, \quad a_j = \left( \frac{M_j}{M_1} \right)^2, \quad a_h = \left( \frac{\mu}{M_1} \right)^2,
\]

the infrared cutoff \( \mu \) is chosen to be equal to 800 GeV due to phenomenological considerations.

So we have

\[
\gamma_{S_j} = 2\gamma_{ij}^{(1)} + 4\gamma_{ij}^{(2)}.
\]

The expressions of reduced cross sections of \( \Delta L = 2 \) RH neutrino exchange processes, which washout a part of the asymmetry are given by

\[
\tilde{\sigma}_{N_j}^{(1)} = \sum_{\alpha=1}^{4} \sum_{\beta=1}^{3} \sum_{j=1}^{4} \left( \lambda_{\alpha j}^* \lambda_{\alpha j} \right) \left( \lambda_{\beta j}^* \lambda_{\beta j} \right) A_{ij}^{(1)} + \sum_{\alpha=1}^{4} \sum_{\beta=1}^{3} \sum_{j=1}^{4} \text{Re} \left( \lambda_{\alpha n}^* \lambda_{\alpha j} \right) \left( \lambda_{\beta n}^* \lambda_{\beta j} \right) B_{nj}^{(1)}
\]

\[
\tilde{\sigma}_{N_j}^{(2)} = \sum_{\alpha=1}^{4} \sum_{\beta=1}^{3} \sum_{j=1}^{4} \left( \lambda_{\alpha j}^* \lambda_{\alpha j} \right) \left( \lambda_{\beta j}^* \lambda_{\beta j} \right) A_{ij}^{(2)} + \sum_{\alpha=1}^{4} \sum_{\beta=1}^{3} \sum_{j=1}^{4} \text{Re} \left( \lambda_{\alpha n}^* \lambda_{\alpha j} \right) \left( \lambda_{\beta n}^* \lambda_{\beta j} \right) B_{nj}^{(2)}
\]
where

\[
A_{ij}^{(1)} = \frac{1}{2\pi} \left[ \frac{1}{D_j} + \frac{a_j}{2D_j^2} - \frac{a_j x}{x + a_j} \left( 1 + \frac{x}{D_j} \right) \ln \left( \frac{x + a_j}{a_j} \right) \right],
\]

\[
A_{ij}^{(2)} = \frac{1}{2\pi} \left[ \frac{x}{x + a_j} + \frac{a_j}{x + 2a_j} \ln \left( \frac{x + a_j}{a_j} \right) \right],
\]

\[
B_{\alpha j}^{(1)} = \frac{\sqrt{a_n a_j}}{2\pi} \left[ \frac{1}{D_j} + \frac{1}{D_n} + \frac{x}{D_j D_n} + \left( 1 + \frac{a_j}{x} \right) \left( \frac{2}{a_n - a_j} - \frac{1}{D_j} \right) \ln \left( \frac{x + a_j}{a_j} \right) \right]
\]

\[+ \left( 1 + \frac{a_n}{x} \right) \left( \frac{2}{a_j - a_n} - \frac{1}{D_j} \right) \ln \left( \frac{x + a_n}{a_n} \right),
\]

\[
B_{\alpha j}^{(2)} = \frac{\sqrt{a_n a_j}}{2\pi} \left[ \frac{x + a_j}{x + a_n + a_j} \ln \left( \frac{x + a_j}{a_j a_n} \right) + \frac{2}{a_n - a_j} \ln \left( \frac{a_n (x + a_j)}{a_j (x + a_n)} \right) \right],
\]

and

\[
D_j = \frac{(x - a_j)^2 + a_j c_j}{x - a_j}, \quad c_j = a_j \sum_{\alpha=1}^{4} \sum_{\beta=1}^{3} \left( \lambda_{\alpha j}^* \lambda_{\alpha j} \lambda_{\beta j}^* \lambda_{\beta j} \right) / 64\pi^2.
\]

then we have

\[
\gamma_W = \sum_{j=1}^{4} \left( \frac{1}{2} \gamma_{D_j} + \frac{Y^2}{\gamma_{N_j}} \gamma_{i_j}^{(1)} + 2 \gamma_{i_j}^{(2)} - \frac{\gamma_{D_j}}{8} \right) + 2\gamma_{N_1}^{(1)} + 2\gamma_{N_1}^{(2)}.
\]

There are several important comments to make at this point. First of all, notice that in eqs. (10), (20) and (21), we have summed over all states \( \ell_\alpha \). However in eqs. (24) and (25) terms which correspond to \( \beta = 4 \), (that is the \( \sigma \) flavour, see figure 11), do not contribute to equation (16) because they are multiplied by \( Y_{\ell_4} = 0 \). Indeed, as was mentioned before, the processes involving at the same time only this flavour in the initial and final state are in thermal equilibrium, and so there is no asymmetry in this flavour. Second, as in the case of FY model, one can parametrize \( \gamma_{D_j} \) and \( \gamma_{\beta j} \) by \( M_1 \) and \( \tilde{m}_{14G} \) (see eqs. (10), (20), (21)). Third, the \( \Delta L = 2 \) processes can be divided and treated in two separate regimes. The contribution in the temperature range from \( z = 1 \) to \( z \approx 10 \) is proportional to \( \sqrt{\tilde{m}_{14G} \tilde{m}_{13G}} \) instead of \( \tilde{m}_{11} \). For the range of temperature \( (z >> 1) \), the dominant contribution at leading order would have been proportional to \( m^2 \equiv tr(m^\dagger \nu m) \), had the summation on the index \( \beta \) gone from 1 to 4. In addition, we emphasize that the so-called RIS (real intermediate states) in the \( \Delta L = 2 \) interactions have to be carefully subtracted, to avoid double counting in the Boltzmann equations. In our calculation, we have followed ref. [13] where the authors have shown how the appropriate \((-1/8)\gamma_{D_j}\) in eq. (31) subtraction must be done.

5 Results and Conclusions

5.1 Constraints on the fourth generation

Our model is based on the addition of a fourth doublet of leptons and quarks (in order to be anomaly-free) to the particle content of the Standard Model, and of four RH neutrino
singlets. The masses and couplings of the singlets are constrained by the left handed neutrino masses (from solar, atmospheric and WMAP data) and the right amount of CP asymmetry $\epsilon_1$. Concerning the fourth generation of leptons and quarks, the direct constraint is that they must be heavier than $M_Z^2$. There are stronger constraints from electroweak precision tests: for example, the lower bound on the charged lepton $\sigma$ from LEP II is approximately 80 GeV \cite{29}. As is well known, the SM central value of the Higgs mass is lower than the direct lower bound set by LEP II \cite{30} and the existence of a fourth generation is one possible way to increase the Higgs mass \cite{31,32,33}. The Z-lineshape versus fourth generation masses have been extensively studied in for example ref. \cite{34} where we can see that the fourth left-handed neutrino is excluded at 95\% CL if its mass is lower than $46.7 \pm 0.2$ GeV. In our analysis the left handed neutrino masses are always constrained in order to fit the data (solar + atmospheric + WMAP) as well as this latter bound for the 4th LH neutrino. There are other constraints on the fourth generation, the generalized CKM matrix elements, etc. which can be found in \cite{26} and in ref. \cite{35}, where present and future experimental searches at Tevatron and LHC are widely discussed.

### 5.2 Numerical solutions of BE

In our numerical analysis we use one of the possible textures for the Yukawa matrix($\lambda_\nu$) in the four generation case which can produce the needed amount of CP-asymmetry and fit the neutrino data, \cite{14}:

$$
\lambda_\nu = C \begin{pmatrix}
\epsilon & \epsilon & \epsilon & \alpha \\
\epsilon & 1 & 1 & 0 \\
\epsilon & 1 & 1 & 0 \\
\epsilon & 0 & 0 & 1/C \\
\end{pmatrix},
$$

where $C, \epsilon$ and $\alpha$ are parameters ($\epsilon$ and $\alpha$ are complex numbers). This texture will induce to first order the following mass matrix for the three lightest LH neutrinos

$$
\mathbf{m_\nu} \propto \begin{pmatrix}
\epsilon^2 & \epsilon & \epsilon \\
\epsilon & 1 & 1 \\
\epsilon & 1 & 1 \\
\end{pmatrix}.
$$

This is a typical texture for a hierarchical spectrum that has to fit the oscillation data (atmospheric, solar, CHOOZ) and the constraints on the absolute mass.
To show the feasibility of our scenario we chose different values of the parameters $C, \epsilon, \alpha$ and values of the RH neutrino masses, $M_i, i = 1, 4$ and we constrain the inputs given the data for the LH neutrino masses and the out-of-equilibrium condition. It is precisely a hierarchy in the Yukawa couplings which allows for an out-of-equilibrium decay without a suppression of the induced CP-asymmetry.

In figure 2 we illustrate for a given set of input parameters the different thermally averaged reaction rates contributing to BE as a function of $z = \frac{M_1}{T}$:

$$\Gamma_X = \frac{\gamma_X}{n_{N_1}} , \quad X = D, S, \Delta L = 2.$$  \hspace{1cm} (34)

It is clear from this plot that for this set of parameters, and this is true for a wide range of parameter space, all rates at $z = 1$ fulfill the out-of-equilibrium condition (i.e. $\Gamma_X < H(z = 1)$), and so the expected washout effect due to the $\Delta L = 2$ processes will be small. The parameters chosen for this illustration are: $M_1 = 450 \text{ GeV}, M_2 = 2 \times 10^6 \text{ GeV}, M_3 = 10^6 \text{ GeV}, M_4 = 605 \text{ GeV}, \epsilon_1 \simeq 4.2 \times 10^{-6}, \tilde{m}_1^{(3G)} = 2.7 \times 10^{-5} \text{ eV}$ and $\tilde{m}_1^{(4G)} = 3.6 \times 10^{-5} \text{ eV}$.

Although a complete scan of the allowed parameter space varying both the values of the RH neutrino masses and the Yukawa couplings has not been done we mention that it is not completely trivial to satisfy all constraints when we vary the parameters. For example, for fixed values of the Yukawa couplings, if we let the second lightest RH neutrino be heavier then the heavier LH neutrino becomes lighter than $M_Z/2$. On the other hand, again for fixed values of the Yukawa couplings, when the masses of the two heaviest RH neutrinos increase or decrease, it becomes difficult to have acceptable light LH neutrino masses. Nevertheless, a more detailed analysis can provide other spectra for the RH neutrinos which are also acceptable.

Figure 3 represents the solution of the BE, abundance and $B-L$ asymmetry, as a function of $z$. The generated value of the baryon asymmetry is $\eta_B \simeq 6 \times 10^{-10}$. Applying the see-saw mechanism to our model for the chosen values of the parameters, we obtain a heavy left-handed neutrino with a mass above 48 GeV, which is consistent with the fourth generation constraints (see section 5.1) and the three light neutrino masses are of the order of $10^{-1} \text{ eV}$ to a few $10^{-7} \text{ eV}$.

We have presented the solutions to the coupled system of Boltzmann equations for our TeV scale model of thermal leptogenesis. We have carefully considered the effect of the interactions involving the heavy fourth generation leptonic fields and consistently written the BE which contribute to the final baryon asymmetry together with the appropriate conversion factor from the lepton asymmetry to the baryonic one. Our results show that in this simple extension of the Standard Model it is possible to produce the right amount of the BAU in a generic way.

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Figure 2: Various thermally averaged reaction rates $\Gamma_X$ contributing to BE normalized to the expansion rate of the Universe $H(z=1)$.

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