Topics in Leptogenesis

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Abstract. Baryogenesis via leptogenesis provides an appealing mechanism to explain the observed baryon asymmetry of the Universe. Recent refinements in the understanding of the dynamics of leptogenesis include detailed studies of the effects of lepton flavors and of the role possibly played by the lepton asymmetries generated in the decays of the heavier singlet neutrinos \( N_2, N_3 \). In this talk I present a short review of these recent developments in the theory of leptogenesis.

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INTRODUCTION

The possibility that the Cosmic Baryon Asymmetry (BAU) could originate from a lepton number asymmetry generated in the \( CP \) violating decays of the heavy seesaw Majorana neutrinos was put forth twenty years ago by Fukugita and Yanagida [1]. Their proposal came shortly after Kuzmin, Rubakov and Shaposhnikov pointed out that above the electroweak phase transition \( B + L \) is violated by fast electroweak anomalous interactions [2]. This implies that any lepton asymmetry would be in part converted into a baryon asymmetry. However, the discovery that at \( T \gtrsim 100 \text{GeV} \) electroweak interactions do not conserve baryon number, also suggested the exciting possibility that baryogenesis could be a purely standard model phenomenon, and opened the way to electroweak baryogenesis [3]. Indeed, in the early 90’s electroweak baryogenesis attracted more interest than leptogenesis. Still, a few remarkable papers appeared that put the first basis for quantitative studies of leptogenesis. Here I will just mention two important contributions: the 1992 Luty’s paper [4] in which the rates for several processes relevant for the Boltzmann equations for leptogenesis were first presented, and the 1996 paper of Covi, Roulet and Vissani [5] that computed the correct expression for the \( CP \) asymmetry in the decays of the lightest Majorana neutrino.

Around year 2000 a flourishing of detailed studies of leptogenesis begins, with a corresponding burst in the number of papers dealing with this subject (see fig. 1). This raise of interest in leptogenesis can be traced back to two main reasons: firstly, by this time it became clear that within the standard model, electroweak baryogenesis fails to reproduce the correct BAU by many orders of magnitude, and that even in supersymmetric models this scenario quite likely is not viable. Secondly, the experimental confirmation (from oscillation experiments) that neutrinos have nonvanishing masses strengthened the

\[ \text{Based on work done in collaboration with G. Engelhard, Y. Grossman, Y. Nir, J. Racker and E. Roulet.} \]
case for the seesaw mechanism, that in turn implies the existence, at some large energy scale, of lepton number violating (\(\mathcal{L}\)) interactions.

The number of important papers and the list of people that contributed to the development of leptogenesis studies and to understand the various implications for the low energy neutrino parameters is too large to be recalled here. Let me just mention the remarkable paper of Giudice et al. [6] that appeared in 2003: in this paper a whole set of thermal corrections for the relevant leptogenesis processes were carefully computed, a couple of mistakes common to previous studies were pointed out and corrected, and a detailed numerical analysis was presented both for the SM and the MSSM cases. Eventually, it was claimed that the residual numerical uncertainties would probably not exceed the 10%-20% level [6]. A couple of years later, Nir, Roulet, Racker and myself [7] carried out a detailed study of additional effects that were not accounted for in the analysis of ref. [6]. This included electroweak and QCD sphaleron effects, the effects of the asymmetry in the Higgs number density, as well as the constraints on the particles asymmetry-densities implied by the spectator reactions that are in thermal equilibrium in the different temperature ranges relevant for leptogenesis [7]. Indeed, we found that the largest of these new effects would barely reach the level of a few tens of percent.

However, two important ingredients that had been overlooked in practically all previous leptogenesis studies, had still to be accounted for. These were the role of the light lepton flavors, and the role of the heavier seesaw Majorana neutrinos. One remarkable exception was the 1999 paper of Barbieri et al. [8] that, besides addressing as the main topic the issue of flavor effects in leptogenesis, pointed out that the lepton number asymmetries generated in the decays of the heavier seesaw neutrinos should also be taken into
account in computing the BAU\textsuperscript{2} However, these important results did not have much impact on subsequent analyses. The reason might be that these were thought to be just order one effects on the final value of the BAU, with no other major consequences for leptogenesis. As I will discuss in the following, the size of the effects could easily reach the level of one order of magnitude, and, most importantly, they can spoil the leptogenesis constraints on the neutrino low energy parameters, and in particular the limit on the absolute scale of neutrino masses \([10]\). This is important, since it was thought that this limit was a firm prediction of leptogenesis with hierarchical seesaw neutrinos, and that the discovery of a neutrino mass \(m_\nu > 0.2\text{eV}\) would have strongly disfavored leptogenesis, or hinted to different scenarios (as e.g. resonant leptogenesis [11]).

**THE STANDARD SCENARIO**

Let us start by writing the first few terms of the leptogenesis Lagrangian:

\[
\mathcal{L} = \frac{1}{2} [\bar{N}_1 (i \not\partial) N_1 - M_1 N_1 N_1] - (\lambda_1 \bar{N}_1 \ell_1 H + \text{h.c.}). \tag{1}
\]

Here \(N_1\) is the lightest right-handed Majorana neutrino with mass \(M_1\), \(H\) is the Higgs field, and \(\ell_1\) is the lepton doublet to which \(N_1\) couples, that when expressed on a complete orthogonal basis \(\{\ell_i\}\) reads

\[
|\ell_1\rangle = (\lambda \lambda^\dagger)^{-1/2} \sum_i \lambda_{1i} |\ell_i\rangle. \tag{2}
\]

In practice it is always convenient to use the basis that diagonalizes the charged lepton Yukawa couplings (the flavor basis) that also has well defined \(CP\) conjugation properties \(CP(\{\ell_i\}) = \{\bar{\ell}_i\}\) with \(i = e, \mu, \tau\). Note that in the first and third term in (1) a lepton number can be assigned to \(N_1\), that is however violated by two units by the mass term. Then eq. (1) implies processes that violate \(L\) like inverse-decays followed by \(N_1\) decays \(\ell_1 \leftrightarrow N_1 \leftrightarrow \bar{\ell}_1\) or off-shell \(\ell_1 H \leftrightarrow \bar{\ell}_1 \bar{H}\) scatterings. At large temperatures \(T \gg M_1\), \(N_1\) decays and inverse decays are blocked because of thermal effects\textsuperscript{3} and off-shell \(L\) processes are suppressed with respect to \(L\)-conserving ones as \(M_1^2/T^2\). At \(T \ll M_1\) decays and inverse decays are Boltzmann suppressed, while \(\ell_1 \leftrightarrow \bar{\ell}_1\) scatterings are suppressed as \(T^2/M_1^2\). Therefore the temperature range in which \(L\) processes can be important for leptogenesis is around \(T \sim M_1\). The possibility of generating an asymmetry between the number of leptons \(n_{\ell_1}\) and antileptons \(n_{\bar{\ell}_1}\) is due to a non-vanishing \(CP\) asymmetry in \(N_1\) decays:

\[
\varepsilon_1 = \frac{\Gamma(N_1 \rightarrow \ell_1 H) - \Gamma(N_1 \rightarrow \bar{\ell}_1 \bar{H})}{\Gamma(N_1 \rightarrow \ell_1 H) + \Gamma(N_1 \rightarrow \bar{\ell}_1 \bar{H})} \neq 0. \tag{3}
\]

\textsuperscript{2} Lepton flavor effects were also considered by Endoh, Morozumi and Xiong in their 2003 paper [9], in the context of the minimal seesaw model with just two right handed neutrinos.

\textsuperscript{3} \(CP\) violating Higgs decays can kinematically occur when \(T \gtrsim 7M_1\) [6] however, because of a large time dilation factor, only a small fraction of the Higgses can actually decay before thermal phase space closes.
In order that a macroscopic $L$ asymmetry can build up, the condition that $L$ reactions are (at least slightly) out of equilibrium at $T \sim M_1$ must also be satisfied. This condition can be expressed in terms of two dimensionful parameters, defined in terms of the Higgs vev $v \equiv \langle H \rangle$ and of the Plank mass $M_P$ as:

$$
\tilde{m}_1 = \frac{(\lambda \lambda^\dagger)_{11} v^2}{M_1}, \quad m_s \approx 10^3 \frac{v^2}{M_P} \approx 10^{-3} \text{eV}.
$$

The first parameter ($\tilde{m}_1$) is related to the rates of $N_1$ processes (like decays and inverse decays) while the second one ($m_s$) is related to the expansion rate of the Universe at $T \sim M_1$. When $\tilde{m}_1 < m_s$ $L$ processes are slower than the expansion and leptogenesis can occur. As $\tilde{m}_1$ increases to values larger than $m_s$, $L$ reactions approach thermal equilibrium thus rendering leptogenesis inefficient because of the back-reactions that tend to erase any macroscopic asymmetry. However, even for $\tilde{m}_1$ as large as $\sim 100 m_s$ a lepton asymmetry sufficient to explain the BAU can be generated. It is customary to refer to the condition $\tilde{m}_1 > m_s$ as to the strong washout regime since washout reactions are rather fast. This regime is considered more likely than the weak washout regime $\tilde{m}_1 < m_s$ in view of the experimental values of the light neutrino mass-squared differences (that are both $> m^2_{\nu}$) together with the theoretical lower bound $\tilde{m}_1 \geq m_{\nu_1}$, where $m_{\nu_1}$ is the mass of the lightest neutrino. The strong washout regime is also theoretically more appealing since the final value of the lepton asymmetry is independent of the particular value of the $N_1$ initial abundance and also of a possible asymmetry $\eta_{\ell_1} = (n_{\ell_1} - n_{\bar{\ell}_1}) / s \neq 0$ ($s$ is the entropy density) preexisting the $N_1$ decay era. This last fact has been often used to argue that for $\tilde{m}_1 > m_s$ only the dynamics of the lightest Majorana neutrino $N_1$ is important, since asymmetries generated in the decays of the heavier $N_{2,3}$ would be efficiently erased by the strong $N_1$-related washouts. As we will see below, the effects of $N_1$ interactions on the $Y_{\ell_{2,3}}$ asymmetries are subtle, and the previous argument is incorrect. The result of numerical integration of the Boltzmann equations for $\eta_{\ell_1}$ can be conveniently expressed in terms of an efficiency factor $\eta_1$, that ranges between 0 and 1:

$$
\eta_1 = 3.9 \times 10^{-3} \eta_1 \varepsilon_1, \quad \eta_1 \approx \frac{m_s}{\tilde{m}_1}.
$$

The second relation gives a rough approximation for $\eta_1$ in the strong washout regime that will become useful in analyzing the impact of flavor effects. The possibility of deriving an upper limit for the the light neutrino masses \cite{10} follows from the existence of a theoretical bound on the maximum value of the $CP$ asymmetry $\varepsilon_1$ (that holds when $N_{1,2,3}$ are sufficiently hierarchical) and relates $M_1$, $m_{\nu_3}$ and the washout parameter $\tilde{m}_1$:

$$
|\varepsilon_1| \leq \left[ \frac{3}{16\pi} \frac{M_1}{v^2} (m_{\nu_3} - m_{\nu_1}) \right] \sqrt{1 - \frac{m^2_{\nu_1}}{\tilde{m}_1^2}}.
$$

The term in square brackets is the so called Davidson-Ibarra limit \cite{12} while the square root is a correction that was first given in \cite{13}. When $m_{\nu_3} \gtrsim 0.1 \text{eV}$, the light neutrinos are quasi-degenerate and $m_{\nu_3} - m_{\nu_1} \approx \Delta m^2_{\text{atm}} / 2m_{\nu_3} \rightarrow 0$ so that, to keep $\varepsilon_1$ finite, $M_1$ is pushed to large values $\gtrsim 10^{15} \text{GeV}$. Since at the same time $\tilde{m}_1$ must remain larger than
m_{\nu_{13}}$ the washout effects also increase, until the surviving lepton-asymmetry is too small to explain the BAU. The interesting limit $m_{\nu_3} \lesssim 0.15$ eV results.

**LEPTON FLAVOR EFFECTS**

In the Lagrangian (1) the terms involving the charged lepton Yukawa couplings have not been included. Since all these couplings are rather small, if leptogenesis occurs at temperatures $T \gtrsim 10^{12}$ GeV, when the Universe is still very young, not many of the related (slow) processes could have occurred during its short lifetime, and leptogenesis has essentially no knowledge of lepton flavors. At $T \lesssim 10^{12}$ GeV the reactions mediated by the tau Yukawa coupling $h_\tau$ become important, and at $T \lesssim 10^9$ GeV also $h_\mu$-reactions have to be accounted for. Including the Yukawa terms for the leptons yields the Lagrangian:

$$\mathcal{L} = \frac{1}{2} [\bar{N}_1 (i \vec{\partial}) N_1 - M_1 N_1 N_1] - (\lambda_{1i} \bar{N}_1 \ell_i H + h_\nu \bar{\ell}_i \ell_i H^\dagger + \text{h.c.}),$$

(7)

where, since we are using the flavor basis, the matrix $h$ of the Yukawa couplings is diagonal. The flavor content of the (anti)lepton doublets $\ell_1$ ($\bar{\ell}_1$) to which $N_1$ decays is now important, since these states do not remain coherent, but are effectively resolved into their flavor components by the fast Yukawa interactions $h_\nu \equiv [8, 14, 15]$. Note that in general, due to CP violating loop effects, $CP(\bar{\ell}_1) \neq \ell_1$, that is the antileptons produced in $N_1$ decays are not the $CP$ conjugate of the leptons, implying that the flavor projections $K_i \equiv |\langle \ell_i | \ell_1 \rangle|^2$ and $\bar{K}_i \equiv |\langle \bar{\ell}_i | \bar{\ell}_1 \rangle|^2$ differ: $\Delta K_i = K_i - \bar{K}_i \neq 0$. The fact that $\ell_1$ and $\bar{\ell}_1$ can differ in their flavor content implies that even when $N_1$ decays with equal branchings into leptons and antileptons (yielding $\epsilon_1 = 0$) the $CP$ asymmetries for the decays into single flavors can still be non-vanishing. The flavor $CP$ asymmetries are defined as $[15]$:

$$\epsilon_1^i = \frac{\Gamma(N_1 \rightarrow \ell_i H) - \Gamma(N_1 \rightarrow \bar{\ell}_i \bar{H})}{\Gamma_{N_1}} = K_i \epsilon_1 + \Delta K_i / 2.$$  

(8)

The factor $\Delta K_i$ in the second equality accounts for the flavor mismatch between leptons and antileptons, while the factor $K_i$ in front of $\epsilon_1$ accounts for the reduction in the strength of the $N_1$-$\ell_1$ coupling with respect to $N_1$-$\bar{\ell}_1$. The same factor also reduces the strength of the washouts for the $i$-flavor, yielding an efficiency factor $\eta_i^1 = \min(\eta_i / K_i, 1)$. Assuming for illustration $\eta_i / K_i < 1$ the resulting asymmetry is

$$Y_L \approx \sum_i \epsilon_1^i \eta_i^1 \approx n_f Y_{\ell_1} + \sum_i \frac{\Delta K_i}{2K_i} \frac{m_+}{m_1}.$$  

(9)

In the first term on the r.h.s. $n_f$ represents the number of flavors effectively resolved by the charged lepton Yukawa interactions ($n_f = 2$ or $3$), while $Y_{\ell_1}$ is the asymmetry that would have been obtained by neglecting the decoherence of $\ell_1$. The second term, that

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$^4$ $\Delta L = 2$ washout processes, that depend on a different parameter than $m_1$, and that can become important when $M_1$ is large, also play a role in establishing the limit.
FIGURE 2. \(|Y_{B-L}|\) (in units of \(10^{-5}|\epsilon_1|\)) as a function of \(K_T\) in two two-flavor regimes. The thick lines correspond to the special case when \(K_T = \bar{K}_T\). The thin lines give an example of the results for \(K_T \neq \bar{K}_T\). The filled circles and squares at \(K_T = 0, 1\) correspond to the aligned cases without flavor effects.

is controlled by the 'flavor mismatch' factor \(\Delta K_i\), can become particularly large in the cases when the flavor \(i\) is almost decoupled from \(N_1\) \((K_i \ll 1)\). This situation is depicted in fig. 2 for the two-flavor case and for two different temperature regimes. The two flat curves give \(|Y_{B-L}|\) as a function of the flavor projector \(K_T\) assuming \(\Delta K_T = 0\), and show rather clearly the enhancement of a factor \(\approx 2\) with respect to the aligned cases (the points at \(K_T = 0, 1\)) for which flavor effects are irrelevant. The other two curves, that peak at values close to the boundaries, when \(\ell_T\) or a combination orthogonal to \(\ell_T\) are almost decoupled from \(N_1\), show that \(\ell_{1,0}^{-1}\) flavor mismatch can produce much larger enhancements. In conclusion, the relevance of flavor effects is at least twofold:

1. The BAU resulting form leptogenesis can be several times larger than what would be obtained neglecting flavor effects.
2. If leptogenesis occurs at temperatures when flavor effects are important, the limit on the light neutrino masses does not hold [14, 16]. This is because there is no analogous of the Davidson-Ibarra bound eq. (6) for the flavor asymmetries \(\epsilon_1^i\).

THE EFFECTS OF THE HEAVIER MAJORANA NEUTRINOS

What about the possible effects of the heavier Majorana neutrinos \(N_{2,3}\) that we have so far neglected? Some recent studies analyzed the so called “\(N_1\)-decoupling” scenario, in which the Yukawa couplings of \(N_1\) are simply too weak to washout an asymmetry generated in \(N_2\) decays (and \(\epsilon_1\) is too small to explain the BAU) [17]. This is a consistent scenario in which \(N_2\) leptogenesis could successfully explain the BAU. However, in the opposite situation when the Yukawa couplings of \(N_1\) are very large, it was generally
assumed that the asymmetries related to $N_{2,3}$ are irrelevant for the computation of the BAU, since they would be washed out during $N_1$ leptogenesis. In contrast to this, in ref. [8] (and more recently also in ref. [18]) it was stated that part of the asymmetry from $N_{2,3}$ decays does in general survive, and must be taken into account when computing the BAU. In ref. [19] Engelhard, Grossman, Nir and myself carried out a detailed study of the fate of a lepton asymmetry preexisting $N_1$ leptogenesis, and we reached conclusions that agree with these statements. I will briefly describe the reasons for this and the importance of the results. Including also $N_{2,3}$ the leptogenesis Lagrangian reads:

$$
\mathcal{L} = \frac{1}{2} [\bar{N}_\alpha (i \not{\partial}) N_\alpha - N_\alpha M_\alpha N_\alpha] - (\lambda_{\alpha i} \bar{N}_\alpha \ell_i H + h.c.),
$$

(10)

where $\alpha = 1, 2, 3$ is a heavy neutrinos index, and the heavy neutrinos are written in the mass basis. It is convenient to define the three (in general non-orthogonal) combinations of lepton doublets $\ell_\alpha$ to which the corresponding $N_\alpha$ decay:

$$
|\ell_\alpha\rangle = (\lambda \lambda^\dagger)_{\alpha\alpha}^{1/2} \sum_i \lambda_{\alpha i} |\ell_i\rangle.
$$

(11)

Let us discuss for definiteness the case where a sizeable asymmetry originates from $N_2$ decays, while $N_1$ leptogenesis is instead rather inefficient. That is, let us assume that $N_2$-related washouts are not too strong, while $N_1$-related washouts are so strong that by itself $N_1$ leptogenesis would not be successful:

$$
\bar{m}_2 \gg m_*, \quad \bar{m}_1 \gg m_*. \quad (12)
$$

To simplify the arguments, let us also impose two additional conditions: thermal leptogenesis, that is a vanishing initial $N_1$ abundance $n_{N_1} (T \gg M_1) \approx 0$, and a strong hierarchy $M_2/M_1 \gg 1$. From this it follows that there are no $N_1$ related washout effects during $N_2$ leptogenesis and, because $n_{N_2}(T \approx M_1)$ is Boltzmann suppressed, there are no $N_2$ related washouts during $N_1$ leptogenesis. Thus $N_2$ and $N_1$ dynamics are decoupled. Now, the second condition in (12) implies that already at $T \gtrsim M_1$ the interactions mediated by the $N_1$ Yukawa couplings are sufficiently fast to quickly destroy the coherence of the state $\ell_2$ produced in $N_2$ decays. Then a statistical mixture of $\ell_1$ and of the states orthogonal to $\ell_1$ builds up, and it can be described by a suitable diagonal density matrix. On general grounds one expects that decoherence effects proceed faster than washout. In the relevant range, $T \gtrsim M_1$, this is also ensured by the fact that because of thermal effects the dominant $\mathcal{O}(\lambda^2)$ washout process (the inverse decay $\ell H \rightarrow N_1$) is blocked [9], and only scatterings with top-quarks and gauge bosons, that have additional suppression factors of $h_t^2$ and $g^2$, contribute to the washout.

Let us consider the case where both $N_2$ and $N_1$ decay at $T \gtrsim 10^{12}$ GeV and flavor effects are irrelevant. In this regime a convenient choice for the orthogonal lepton basis is $(\ell_1, \ell_0, \ell'_0)$ where, without loss of generality, $\ell'_0$ can be chosen to satisfy $\langle \ell'_0 | \ell_2 \rangle = 0$. Then the asymmetry $Y_{\ell_2}$ produced in $N_2$ decays decomposes in the two components:

$$
Y_{\ell_0} = c^2 Y_{\ell_2}, \quad Y_{\ell_1} = s^2 Y_{\ell_2},
$$

(13)

where $c^2 \equiv |\langle \ell_0 | \ell_2 \rangle|^2$ and $s^2 = 1 - c^2$. The crucial point here is that in general we expect $c^2 \neq 0$, and since $\ell_0$ is orthogonal to $\ell_1$, $Y_{\ell_0}$ is protected against $N_1$ washouts. Then a
finite part of the asymmetry $Y_{\ell_2}$ from $N_2$ decays survives through $N_1$ leptogenesis. A more detailed study \[19\] reveals also some unexpected features. For example, $Y_{\ell_1}$ is not driven to zero in spite of the strong $N_1$-related washouts, rather, only the sum of $Y_{\ell_1}$ and of the Higgs asymmetry $Y_H$ vanishes, but not the two separately. (This can be traced back to the presence of a conserved charge related to $Y_{\ell_0}$.) As a result we obtain $Y_{\ell_1} = -Y_{\ell_0}/4$ and $Y_H = Y_{\ell_0}/2$, while the total lepton asymmetry is $Y_L = (3/2)Y_{\ell_0} = (3/2)c^2Y_{\ell_2}$.

For $10^9 \lesssim M_1 \lesssim 10^{12}$ GeV the $\ell_{2,1}$ flavor structures are only partially resolved during $N_1$ leptogenesis, and a similar result is obtained. However, when $M_1 \lesssim 10^9$ GeV and the full flavor basis ($\ell_e, \ell_\mu, \ell_\tau$) is resolved, there are no directions in flavor space where an asymmetry can remain protected, and then the whole $Y_{\ell_2}$ can be erased. In conclusion, the common assumption that when $N_1$ leptogenesis occurs in the strong washout regime the final BAU is independent of initial conditions, does not hold in general, and is justified only in the following cases \[19\]: i) Vanishing decay asymmetries and/or efficiency factors for $N_{2,3}$ ($\epsilon_2 \eta_2 \approx 0$ and $\epsilon_3 \eta_3 \approx 0$); ii) $N_1$-related washouts are still significant at $T \lesssim 10^9$ GeV; iii) Reheating occurs at a temperature in between $M_2$ and $M_1$. In all other cases the $N_{2,3}$-related parameters cannot be ignored when calculating the BAU, and any constraint inferred from analyses based only on $N_1$ leptogenesis are not reliable.

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