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Strong thermal leptogenesis and the N_2-dominated scenario

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Abstract

We briefly review the main aspects of leptogenesis, pointing out the main reasons that draw attention to the so-called N_2-dominated scenario. We consider the conditions that stem out when the final asymmetry is required to be fully independent of the initial conditions. We show that in this scenario, called strong thermal leptogenesis, when barring special cancellations in the seesaw formula and in the flavoured decay parameters, a lightest neutrino mass $m_1 \gtrsim 10$ meV for normal ordering and $m_1 \gtrsim 3$ meV for inverted ordering are favoured. We then focus on the $SO(10)$-inspired leptogenesis models that naturally realise N_2-dominated leptogenesis. We show how the combination with the strong thermal leptogenesis conditions yields important predictions on neutrino parameters. Finally, we briefly comment on the power of forthcoming neutrino experiments to either support or severely corner these leptogenesis scenarios.

1. Introduction

Leptogenesis is a particularly attractive way for producing the baryon asymmetry of the Universe, since it relies on the seesaw mechanism that is able to explain the observed neutrino masses and mixing.

We focus here on the standard type-I seesaw, with the inclusion of heavy right-handed (RH) Majorana neutrinos, N_{Ri}, that couple to the lepton doublets via Yukawa interactions. When the RH neutrinos mass scale is much higher than the electroweak scale (the so called seesaw limit), the neutrino mass spectrum splits into two set. One is made of light, Majorana, active neutrinos ν_i whose mass matrix is given by

$$m_\nu = -m_D\, D_M^{-1}\, m_D^T. \tag{1}$$

Here D_M is the diagonal Majorana mass matrix and $m_D \equiv v y$ is the Dirac neutrino mass matrix, proportional to the Higgs vacuum expectation value v and the RH neutrinos Yukawa couplings y. The eigenvalues of m_ν can be obtained by means of the PMNS matrix $D_m = -U^\dagger\, m_\nu\, U^*$, and turn out to be suitably small, without the need for artificially tuned Yukawa couplings.
The second set is composed of heavy Majorana neutrinos N_i, whose mass matrix basically coincides with D_M up to negligible corrections.

This elegant way of generating small neutrino masses has however some drawbacks. The addition of heavy RH neutrinos introduces new free parameters to the model. For instance, if we assume, as we will always do in what follows, the presence of 3 RH neutrinos, the seesaw model depends on 18 free parameters that split into the 9 so-called *low-energy neutrino parameters,* probed in experiments, and a further 9 describing the high-energy scale of the heavy neutrinos. The seesaw mechanism on its own does not provide any explanation nor prediction on their values. In this regard, we can then either seek further theoretical input, or turn to an apparently unrelated phenomenology such as the baryon asymmetry of the Universe. This can be done by means of leptogenesis.

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2. Leptogenesis and the N_2-dominated scenario

Leptogenesis [1] can be realised within a standard type-I seesaw model with the inclusion of at least 2 RH neutrinos. As already mentioned, in this work we shall consider 3 RH neutrinos. Their Majorana nature automatically ensures lepton-number violation. The presence in the seesaw lagrangian of Yukawa couplings to the LH lepton doublets allows the heavy neutrinos, through their RH component, to decay into leptons and antileptons. Due to complex Yukawa couplings and the interference between tree-level and one-loop diagrams, the heavy neutrinos decays are CP-violating.

CP-violation is estimated by the CP-asymmetry factors

$$\varepsilon_{i\alpha} \equiv -\frac{\Gamma_{i\alpha} - \bar{\Gamma}_{i\alpha}}{\sum_\beta \left(\Gamma_{i\beta} + \bar{\Gamma}_{i\beta}\right)}, \tag{2}$$

where $\Gamma_{i\alpha}$ ($\bar{\Gamma}_{i\alpha}$) are the zero-temperature decay rates of the heavy neutrino N_i into leptons (antileptons) of flavour α. Decays are also regulated by the so-called flavoured decay factors

$$K_{i\alpha} \equiv \frac{\Gamma_{i\alpha}}{H(T = M_i)}, \tag{3}$$

where $H(T = M_i)$ is the value of the Hubble parameter when N_i becomes non-relativistic. The decay factors $K_{i\alpha}$ do not only measure whether decays are in equilibrium at that epoch but also give an estimate of the efficiency of the inverse decays, responsible for the production of N_i. Inverse decays are also the main source of washout of the generated asymmetry, therefore $K_{i\alpha}$ are a measure of the washout strength too.

When occurring out of equilibrium, heavy neutrino decays can produce a net lepton asymmetry. If this takes place at temperature $T \gtrsim 100$ GeV, electroweak sphaleron processes are able to partly convert the lepton asymmetry into the baryon sector.

We focus, now, on a particular hierarchical heavy neutrino spectrum with

$$10^{12}\,\text{GeV} \lesssim M_3, \tag{4}$$

$$10^{9}\,\text{GeV} \lesssim M_2 \lesssim 10^{12}\,\text{GeV}, \tag{5}$$

$$M_1 \lesssim 10^{9}\,\text{GeV}. \tag{6}$$

Since N_3 is the heaviest neutrino, its contribution to the asymmetry is negligible, and so is that of N_1, which is too light. Therefore, the final asymmetry is mainly produced by the next-to-lightest heavy neutrino and this scenario takes the name N_2*-dominated* [2]. If flavour interactions are completely neglected, heavy neutrinos decay into a coherent superposition of flavour eigenstates. The final value of the asymmetry is exponentially suppressed by N_1's washout. Only when $K_1 \ll 1$ the model can produce the observed asymmetry. A numerical analysis shows that this is realised in only $\sim 0.2\%$ of the parameter space [3].

This picture improves by taking account of flavour interactions. Considering that τ-interactions are efficient for $T \lesssim 10^{12}$ GeV, and muonic ones for $T \lesssim 10^9$ GeV, and given the heavy neutrino spectrum as in eqs. (4-6), N_2 decays in a two fully-flavoured regime, while the washout due to N_1 takes place in a three fully-flavoured regime. Around the transition temperatures, the behaviour cannot be described by the usual Boltzmann equations and a density matrix formalism must be adopted [4]. Avoiding these regions, N_1's washout acts separately on each flavour $\alpha = e, \mu, \tau$, exponentially suppressing the asymmetry by the relative flavoured decay parameter $K_{1\alpha}$ [5, 6, 7, 4]. Therefore, it is sufficient that only one $K_{1\alpha}$ is smaller than 1 in order to have enough asymmetry produced. This is found to happen in $\sim 30\%$ of the parameter space, showing that thanks to flavour effects the N_2-dominated scenario can represent a viable model of leptogenesis.

3. Strong thermal leptogenesis

The final baryon asymmetry depends in general on the initial conditions. Moreover, the high reheating temperatures $T_{\rm RH} \gtrsim 10^{12}$ GeV required by this scenario of thermal leptogenesis[1], see eq. (5), allow several mechanism (e.g. GUT-, gravitational-, Affleck-Dine-baryogenesis) to produce a rather high asymmetry. This asymmetry is produced before the onset of leptogenesis, and is referred to as *initial pre-existing asymmetry*, $N^{\rm p,i}_{B-L}$. At lower temperatures, leptogenesis takes place and in general it modifies this asymmetry to a final value $N^{\rm p,f}_{B-L}$. We can require

$$N^{\rm p,f}_{B-L} \ll N^{\rm lep,f}_{B-L} \simeq N^{\rm CMB}_{B-L}, \tag{7}$$

namely that the final pre-existing asymmetry is negligible in comparison to the asymmetry produced by leptogenesis, $N^{\rm lep,f}_{B-L}$, which therefore genuinely constitutes the amount of asymmetry measured today. In this way, full independence of the initial conditions is ensured and leptogenesis is said to satisfy the *strong thermal* condition [9]. With hierarchical heavy neutrino spectra, it

[1]In this respect, it can be interesting to notice that these high values of $T_{\rm RH}$ would be even quite natural if the BICEP2 claim [8] is confirmed.

has been shown [9] that strong thermal leptogenesis can be realised, but only in a two-stage process where the heavy neutrinos show a spectrum as in eqs. (4-6). This is one of the main reasons why the N_2-dominated scenario is important.

In this scenario, N_2 decays in the two-flavoured regime and efficiently washes out the τ component of $N^{p,i}_{B-L}$. On the other hand, N_1 decays in the three-flavoured regime, separately erasing the e and μ components of $N^{p,i}_{B-L}$, while the produced asymmetry survives in the τ flavour. This translates into a set of conditions on the relevant decay parameters

$$K_{1e}, K_{1\mu}, K_{2\tau} \gg 1, \quad K_{1\tau} \lesssim 1. \tag{8}$$

When these conditions are imposed on the requirement of successful leptogenesis (i.e. $\eta_B^{\rm lep} \simeq \eta_B^{\rm CMB}$) a precise analytical lower bound on the absolute mass scale m_1 appears [10]. Adopting the Casas-Ibarra parameterisation [11], we introduce a complex orthogonal matrix $\Omega = D_m^{-1/2}\, U^\dagger V_L^\dagger\, D_{m_D}\, U_R D_M^{-1/2}$, that links the low-energy (neutrino masses and mixing) and the high-energy (heavy neutrinos) parameters. In this way, the expression for the decay parameters becomes

$$K_{i\alpha} = \left| \sum_j \sqrt{\frac{m_j}{m_*}}\, U_{\alpha j}\, \Omega_{ji} \right|^2 . \tag{9}$$

Considering low values of m_1, that is $m_1 \lesssim m_{\rm sol}$, we can simplify the expression for $K_{1\alpha}$

$$K_{1\alpha} \simeq \left| \sqrt{\frac{m_1}{m_*}} U_{\alpha 1}\, \Omega_{11} + \sqrt{\frac{m_{\rm sol}}{m_*}} U_{\alpha 2}\, \Omega_{21} + \sqrt{\frac{m_{\rm atm}}{m_*}} U_{\alpha 3}\, \Omega_{31} \right|^2 . \tag{10}$$

Specifying $\alpha = \tau$, we obtain

$$\sqrt{\frac{m_{\rm atm}}{m_*}} U_{\tau 3}\, \Omega_{31} = -\sqrt{\frac{m_1}{m_*}} U_{\tau 1}\, \Omega_{11} - \sqrt{\frac{m_{\rm sol}}{m_*}} U_{\tau 2}\, \Omega_{21} + \sqrt{K_{1\tau}}\, e^{i\varphi_\tau}, \tag{11}$$

and substitute it in the expression for $K_{2\delta}$, with $\delta = e, \mu$

$$K_{1\delta} = \left| \Omega_{11} \sqrt{\frac{m_1}{m_*}} \left(U_{\delta 1} - \frac{U_{\tau 1}}{U_{\tau 3}} U_{\delta 3} \right) + \sqrt{K^0_{1\delta}} e^{i\varphi_0} \right|^2 . \tag{12}$$

Here $K^0_{1\delta} \equiv K_{1\delta}(m_1 = 0)$, so that

$$\sqrt{K^0_{1\delta}}\, e^{i\varphi_0} \equiv \Omega_{21} \sqrt{\frac{m_{\rm sol}}{m_*}} \left(U_{\delta 2} - \frac{U_{\tau 2}}{U_{\tau 3}} \right) + \frac{U_{\delta 3}}{U_{\tau 3}} \sqrt{K_{1\tau}} e^{i\varphi_\tau}. \tag{13}$$

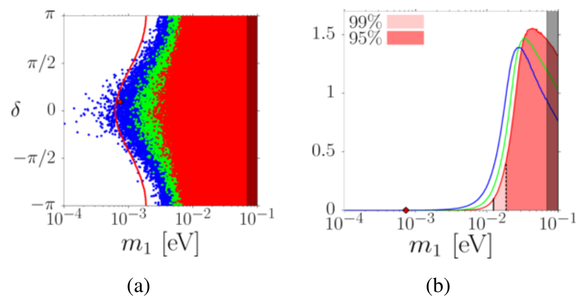


Figure 1: (a) Scatter plot of successful strong leptogenesis models. Initial pre-existing asymmetries are $N^{\rm p,i}_{B-L} = 0.1, 0.01, 0.001$ respectively in red, green and blue. The solid line is the 95% confidence level analytical lower bound. (b) Distribution of probability of m_1 from the scatter plot, colour code as in (a). These models have $M_\Omega = 2$ [10].

We can take the size of the entries of the orthogonal matrix to be in general $|\Omega_{ij}|^2 \le M_\Omega$. Usually $M_\Omega = \mathcal{O}(1)$, so that the seesaw mechanism truly relies on the interplay between the electroweak and the Majorana scales, rather than on fine-tuned cancellations. This way, the maximum values of $K_{1\delta}$ are obtained for

$$K^{\rm max}_{1\delta} = \left(M_\Omega \sqrt{\frac{m_1}{m_*}} \left| U_{\delta 1} - \frac{U_{\tau 1}}{U_{\tau 3}} U_{\delta 3} \right| + \sqrt{K^{0,\rm max}_{1\delta}} \right)^2 , \tag{14}$$

with

$$\sqrt{K^{0,\rm max}_{1\delta}} \equiv M_\Omega \sqrt{\frac{m_{\rm sol}}{m_*}} \left| U_{\delta 2} - \frac{U_{\tau 2}}{U_{\tau 3}} \right| + \left| \frac{U_{\delta 3}}{U_{\tau 3}} \right| \sqrt{K^{\rm max}_{1\tau}}. \tag{15}$$

The strong thermal leptogenesis conditions in eq. (8) can be specified by using the general relation in eq. (7). For each flavour we must have $N^{\rm p,f}_{\Delta\alpha} < \zeta N^{\rm lep,f}_{\Delta\alpha}$, with $\zeta \ll 1$, where $N^{\rm p,i}_{\Delta\alpha}$ is the fraction of pre-existing asymmetry in the flavour α. Since the washout takes place exponentially, we can define

$$K_{\rm st}\left(N^{\rm p,i}_{\Delta\alpha}\right) \equiv \frac{3}{8\pi} \left[\ln\left(\frac{100\zeta}{N^{\rm CMB}_{B-L}} \right) + \ln |N^{\rm p,i}_{\Delta\alpha}| \right] \tag{16}$$

$$\simeq 18 + 0.85 \left(\ln\zeta + \ln |N^{\rm p,i}_{\Delta\delta}| \right), \tag{17}$$

and so we must require for each flavour $\delta = e, \mu$

$$K_{1\delta} \gtrsim K_{\rm st}. \tag{18}$$

This in turn implies $K^{\rm max}_{1\delta} > K_{\rm st}(N^{\rm p,i}_{\Delta\delta})$, namely

$$\left(M_\Omega \sqrt{\frac{m_1}{m_*}} \left| U_{\delta 1} - \frac{U_{\tau 1}}{U_{\tau 3}} U_{\delta 3} \right| + \sqrt{K^{0,\rm max}_{1\delta}} \right)^2 > K_{\rm st}(N^{\rm p,i}_{\Delta\delta}). \tag{19}$$

Solving with respect to m_1, we finally obtain

$$m_1 > \frac{m_*}{M_\Omega} \max_\delta \left[\left(\frac{\sqrt{K_{\rm st}(N^{\rm p,i}_{\Delta\delta})} - \sqrt{K^{0,\rm max}_{1\delta}}}{\left| U_{\delta 1} - \frac{U_{\tau 1}}{U_{\tau 3}} U_{\delta 3} \right|} \right)^2 \right], \quad (20)$$

when

$$K^{0,\rm max}_{1\delta} < K_{\rm st}(N^{\rm p,i}_{\Delta\delta}). \quad (21)$$

For normal ordering (NO), the maximum value is obtained for $\delta = e$. Moreover, imposing the other strong thermal condition, i.e. $K_{1\tau} \lesssim 1$, the condition in eq. (21) can be satisfied and a lower bound on m_1 can actually be found. Considering $M_\Omega = 2$, $N^{\rm p,i}_{B-L} = 0.1$, taking the experimental values of the mixing angles and $\delta = 0$ we have $m_1^{\rm lb} \simeq 0.7\,{\rm meV}$. This analytical lower bound holds similarly for inverted ordering (IO), but in this case it is generally much looser, due to the replacement $m_{\rm sol} \to m_{\rm atm}$. More precisely, an analytical lower bound on m_1 holds for IO only when $M_\Omega \lesssim 0.9$ and is therefore much looser than for NO.

From a statistical analysis of the simulations, we have noticed that the actual analytical lower bound is hardly saturated. On the contrary, strong thermal leptogenesis models tend to prefer rather high values of the absolute neutrino mass scale. For natural values of M_Ω (i.e. $M_\Omega \leq 4$), around 99% of models have $m_1 \gtrsim 10\,{\rm meV}$ in NO, while for IO the statistical limit is looser and 99% of models show $m_1 \gtrsim 3\,{\rm meV}$.

This result is particularly interesting since it provides us with a precise prediction that makes leptogenesis, in its strong thermal version, testable at future experiments. It can face many evidences from cosmological observations, hwowever, the sensitivity is not yet sufficient to reasonably discriminate from the hierarchical ($m_1 = 0$) limit. This scenario will improve in forthcoming experiments, if a precision $\delta(\sum m_{\nu_i}) \lesssim 10\,{\rm meV}$ is reached [12]. In this case, measurements of the sum of the neutrino masses will be able to either support or severely corner strong thermal leptogenesis. It must be noticed that it is particularly important that long-baseline neutrino oscillation experiments determine the ordering. As we have seen, NO and IO show rather different features and, in particular, the lower bound for IO is much looser.

It can be convenient to point out once again how the simple requirement of full independence of the initial conditions can lead us to interesting constraints on the absolute neutrino mass scale, thus putting strong thermal leptogenesis within the experimental reach.

3.1. *SO*(10)*-inspired leptogenesis*

As already mentioned, the attractive features of the seesaw mechanism rely on the addition of extra particles. These particles are somehow added by hand to the SM lagrangian and so is their high mass scale. A more elegant and attractive origin of the RH neutrinos and their mass scale can be found in Grand Unified Theories (GUT). In particular, it can be noticed that models based on $SO(10)$ as unification group [13, 14, 15, 16], naturally include three RH neutrinos in the same irreducible representation together with quarks and leptons. RH neutrinos precisely fit in the $\overline{16}$ representation of $SO(10)$ and their mass scale is linked to low-energy parameters, such as the Dirac neutrino mass matrix m_D and the PMNS mixing matrix U. With this in mind, it can be interesting to apply some conditions, that are directly inspired to $SO(10)$ GUT models, to leptogenesis. In particular:

- diagonalising the Dirac mass matrix as $m_D = V_L^\dagger D_{m_D} U_R$, the diagonal matrix is set to
$$D_{m_D} = {\rm diag}\,(\alpha_1 m_u,\ \alpha_2 m_c,\ \alpha_3 m_t)\,, \quad (22)$$
where m_u, m_c, m_t are the up-quark masses at the leptogenesis scale,
- the proportionality constants are fixed to be $0.1 \lesssim \alpha_i \lesssim 10$,
- the unitary matrix V_L is set to be $\mathbb{1} \leq V_L \leq V_{\rm CKM}$, where $V_{\rm CKM}$ is the Cabibbo-Kobayashi-Maskawa quark mixing matrix.

By imposing these relations, together with the strong thermal leptogenesis conditions in eq. (8), several precise predictions on the heavy neutrino mass spectrum and the low-energy neutrino parameters can be obtained. Numerical simulations [17] have pointed out the following list of results.

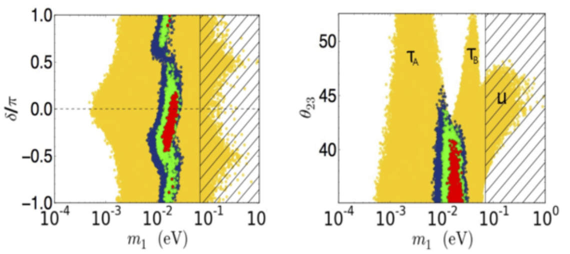
i) The ordering of the active neutrino spectrum must be normal,

ii) the absolute neutrino mass scale is $m_1 \simeq 20\,{\rm meV}$,

iii) the neutrinoless double-beta decay ($0\nu\beta\beta$) effective neutrino mass is found to be $|m_{\nu ee}| \simeq 0.8\,m_1 \simeq 15\,{\rm meV}$,

iv)[2] a lower bound on the reactor mixing angle is found $\theta_{13} \gtrsim 2°(0.5°)$ for $N^{\rm p,i}_{B-L} = 10^{-1}(10^{-2})$,

[2]This lower bound was found in [18], before the measurement by DayaBay [19].



(a) (b)

Figure 2: Scatter plot of successful $SO(10)$-inspired, strong thermal leptogenesis models in the plane (m_1, δ), fig. (a), and (m_1, θ_{23}) in (b). Yellow, blue, green and red dots correspond to $N^{\rm p,i}_{B-L} = 0,\ 10^{-3},\ 10^{-2},\ 10^{-1}$. Notice the predicted values of m_1, δ and θ_{23} for $N^{\rm p,i}_{B-L} = 0.1$ [17].

v) an upper bound on the atmospheric mixing angle is found $\theta_{23} \lesssim 41^\circ (43^\circ)$ for $N^{\rm p,i}_{B-L} = 10^{-1} (10^{-2})$,

vi) the Dirac CP-violating phase must take values $\delta \in [-\pi/2, \pi/5]$ for $N^{\rm p,i}_{B-L} = 10^{-1}$, while for $N^{\rm p,i}_{B-L} = 10^{-2}$ we have $\delta \notin [0.4\pi, 0.7\pi]$,

vii) the Majorana phases tend to cluster around $(\sigma, \rho) \simeq (0.8 + n,\ 1.25 + n)\pi$ or $(\sigma, \rho) \simeq (0.7 + n,\ 0.75 + n)\pi$, with $n = 0, 1$.

In general, we can notice that prediction (ii) is consistent with the lower bound in eq. (20) from strong thermal leptogenesis.

Most importantly, $SO(10)$ is responsible for the upper bound on θ_{23}, that is then constrained to the value in (v) by strong thermal conditions. This prediction is particularly interesting, together with the value of δ given by (vi), and are shown in fig. 2. They both provide us with definite values that can be tested at the experiments. The value in (vi) is intriguingly in very good agreement with the current central value of the neutrino global fits [20, 21] and will be tested in long-baseline neutrino oscillation experiments (e.g. NOνA). On the other hand, the octant of θ_{23} is still very unstable.

4. Conclusions

We have shown that the theoretical request of full independence of initial condition is able to constrain the parameters space and give interesting results on neutrino masses and mixing. This goes in the direction of the need for "testable" leptogenesis models. The lack of strong signals of new Physics, except in Cosmology and in the neutrino sector, can push physical speculation towards these research fields, but it must not be forgotten that each proposed model must avoid the problem of difficult testability. With this in mind, it is evident that the restriction to a minimal scenario of leptogenesis, like that provided by type-I seesaw, is not enough to get solid predictions, even when asking for successful leptogenesis. This can be avoided when strong thermal leptogenesis is required, on its own, or in combination with conditions inspired to $SO(10)$ GUT theories. Both these conditions focus the attention on the so-called N_2-dominated scenario, with the heavy neutrino spectrum as in eqs. (4-6). This scenario is often neglected when dealing with leptogenesis, because one may expect that the washout by the lightest neutrino totally erases the lepton asymmetry generated by the next-to-lightest. However, it can be shown that this is not true when flavour effects are considered, so that the asymmetry can survive N_1's washout in some flavour directions, along which the washout is mild. This must not be regarded as fine-tuned, since successful leptogenesis can be produced by a flavoured N_2-dominated scenario in around 30% of the parameter space [3].

Moreover, in this scenario, favourable values of the absolute neutrino mass scale are found for NO, $m_1 \gtrsim 10\,\mathrm{meV}$ and IO, $m_1 \gtrsim 3$ meV. In NO a further analytical lower bound is found for natural choices of Ω ($M_\Omega \lesssim 4$), see eq. (20), while in IO the analytical threshold is obtained only for $M_\Omega \lesssim 0.9$. These constraints on m_1 allow the future cosmological observations, together with the determination of the neutrino mass ordering, to test strong thermal leptogenesis.

When $SO(10)$-inspired conditions are taken into account, a much richer landscape of predictions arises. In particular, negative values of the Dirac CP-violating phase, $\delta \simeq -\pi/4$, and first-octant atmospheric mixing angle, $\theta_{23} \lesssim 41^\circ$, are predicted. This will enable forthcoming neutrino oscillation experiments to seriously put to test such scenarios of thermal leptogenesis.

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