LEPTOGENESIS*

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Leptogenesis is a cosmological consequence of the seesaw mechanism and provides a link between the observed baryon asymmetry and neutrino masses. We show which information can be inferred from leptogenesis on neutrino masses and on those high energy seesaw parameters that represent a sort of 'dark side' for conventional experiments. We also report on a new scenario of leptogenesis that opens new opportunities for leptogenesis to be tested.

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1. Introduction

Current cosmological observations support the idea that the observed matter–anti-matter asymmetry of the Universe is the relic trace of a dynamical generation process [1], so called baryogenesis, occurred during the very early Universe history. The asymmetry is measured in the form of a *baryon to photon number ratio* at the time of recombination. A combination of WMAP and SLOAN data gives [2]

$$\eta_B^{\rm CMB} = (6.3 \pm 0.3) \times 10^{-10} \,. \tag{1}$$

Even though all three Sakharov's conditions for successful baryogenesis are satisfied in the Standard Model, the final predicted asymmetry would fall by far below the observed value. Therefore, an explanation of the observed asymmetry requires 'new physics'. The discovery of neutrino masses is a strong indication of physics beyond the Standard Model. It is thus remarkable that the most elegant way to understand neutrino masses, the seesaw mechanism, can also provide a simple explanation of the matter–anti-matter

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asymmetry of the Universe known as Leptogenesis [3]. Conversely Leptogenesis can be also regarded as a powerful cosmological tool to test the seesaw mechanism, together with the other more conventional phenomenologies, such as neutrino mixing and, hopefully to be observed in the next years, $\beta\beta 0\nu$ decay and CP violation in neutrino mixing. Here we will concentrate our attention on leptogenesis, showing the information that can be derived on the seesaw parameters.

2. From seesaw to leptogenesis

Adding to the Standard Model Lagrangian three right handed (RH) neutrinos with Yukawa couplings h and a Majorana mass term M, after spontaneous symmetry breaking, a Dirac neutrino mass term, $m_{\rm D} = h\,{\rm v}$, is generated by the vacuum expectation value of the Higgs boson and the whole neutrino mass term can be written [1] as

$$\mathcal{L}^{\nu}_{\rm mass} = -\frac{1}{2}\left[(\bar{\nu}_{\rm L}^{c}, \bar{\nu}_{\rm R}) \left(\begin{array}{cc} 0 & m_{\rm D}^{T} \\ m_{\rm D} & M_{\rm R} \end{array} \right) \left(\begin{array}{c} \nu_{\rm L} \\ \nu_{\rm R}^{c} \end{array} \right)\right] + \text{h.c.} \tag{2}$$

Assuming that the eigenvalues of M are much higher than those of $m_{\rm D}$, then one has a splitting between 3 heavy eigenstates with masses $M_1 \leq M_2 \leq M_3$, approximately coinciding with the eigenvalues of M, and 3 light eigenstates with masses $m_1 \leq m_2 \leq m_3$ given by the seesaw formula

$$m_{\nu} = -m_{\rm D}\,\frac{1}{M}\,m_{\rm D}^{T}\,. \tag{3}$$

All matrices are in general complex and this provides a natural source of CP violation, one of the three Sakharov necessary conditions for successful baryogenesis. Indeed, considering the decays of the RH neutrinos N_i, one has that these can proceed with a different rate into leptons ($N_i \rightarrow l\,\phi^{\dagger}$) and anti-leptons ($N_i \rightarrow \bar{l}\,\phi$). The difference can be expressed in terms of the CP asymmetry parameter defined as

$$\varepsilon_i \equiv -\frac{(\Gamma_i - \bar{\Gamma}_i)}{(\Gamma_i + \bar{\Gamma}_i)}\,. \tag{4}$$

If the mass differences $|M_i - M_j|$ are not too small compared to the rate differences, then ε_i can be calculated from the interference between tree level and one loop graphs, self energy plus vertex correction [4]. It has been pointed out [5] that, working in a basis where the Majorana mass term is diagonal, such that $M \rightarrow D_M \equiv \text{diag}(M_1, M_2, M_3)$, the seesaw formula is

[1] We consider the case where a triplet Higgs term is absent.

equivalent to an orthogonality condition for a matrix Ω through which the Dirac mass matrix can be parameterized as

$$m_{\mathrm{D}} = U\, D_m^{1/2}\, \Omega\; D_M^{1/2}\,. \tag{5}$$

Here U is the matrix that diagonalizes m_ν, such that

$$D_m \equiv \mathrm{diag}(m_1, m_2, m_3) = -U^\star\, m_\nu\, U^\dagger\,. \tag{6}$$

It can be identified with the MNS neutrino mixing matrix in a basis where charged leptons are diagonal. This parametrization is particularly useful for leptogenesis. First of all it shows that there are 18 parameters: 9 'low energy parameters' (the 3 light neutrino masses m_i and the 6 parameters in the MNS mixing matrix U) and 9 'high energy parameters' (the 3 RH neutrino masses M_i and the 6 parameters in the orthogonal matrix Ω). From neutrino mixing experiments we have information on some of the parameters in U, in particular we measure two mass squared differences:

$$(m_3^2 - m_1^2)^{1/2} = m_{\mathrm{atm}} \simeq 0.05\,\mathrm{eV} \tag{7}$$

and

$$(m_{2(3)}^2 - m_{1(2)}^2)^{1/2} = m_{\mathrm{sol}} \simeq 0.009\,\mathrm{eV} \tag{8}$$

for a normal (inverted) scheme. We still miss a determination of the *absolute neutrino mass scale*, that we will indicate in terms of the lightest neutrino mass m_1. The CP asymmetries ε_i depend only on the $m_{\mathrm{D}}^\dagger\, m_{\mathrm{D}}$ entries, and from (5) one can see that the mixing matrix U cancels out, implying that the $\not\!\!\mathrm{CP}$ responsible for leptogenesis stems uniquely from Ω and is *in general* not dependent on U. Vice versa the high energy parameters do not enter the neutrino mass matrix m_ν and thus $\not\!\!\mathrm{CP}$ in neutrino mixing can only arise from U. Thus, the orthogonal parametrization shows a full disentanglement of $\not\!\!\mathrm{CP}$ in neutrino mixing and in leptogenesis. An important consequence is that it is not possible to prove or disprove leptogenesis by measuring or constraining $\not\!\!\mathrm{CP}$ in neutrino mixing.

On the other hand many possible connections between the leptogenesis CP asymmetry and $\not\!\!\mathrm{CP}$ in neutrino mixing have been worked out within specific frameworks: this is not in contradiction with the previous conclusion. Indeed suppose one adds, to the general framework, some extra theoretical input or phenomenological information that specifies or measures a subset $\boldsymbol{X}$ of the seesaw parameters in terms of which one can replace part or all the Ω parameters through a non trivial transformation $\Omega = \Omega(m_i, M_i, U; \boldsymbol{X})$. This brings to a new parametrization $(M_i, m_i, U, \boldsymbol{X})$, where now the $\boldsymbol{X}$'s have to be regarded as known parameters. In this way one has a reduction of the number of parameters that allows to express Ω as a (non trivial) function $\Omega(m_i, M_i, U)$. When this is plugged into the ε_i's, the CP asymmetry responsible for leptogenesis becomes a function of U. Thus a measurement of $\not{C}\!P$ in neutrino mixing, more than a test of leptogenesis, will provide a test of the specific '$\boldsymbol{X}$-framework'. However, *common sense* suggests that 'if there are phases in U, then why not also in Ω' and so the existence of such a framework is quite reasonable, Therefore, a detection of $\not{C}\!P$ in neutrino mixing, if not a smoking gun, will certainly provide an additional piece in support of leptogenesis.

The orthogonal parametrization (*cf.* (5)) is also a useful technical tool to solve different problems in leptogenesis. The final asymmetry can be in general written as the sum of the contributions from the decays of all three N_i,

$$\eta_B = d \sum_i \varepsilon_i \, \kappa_i \, , \tag{9}$$

where $d = a_{\rm sph}/f \simeq 10^{-2}$ takes into account both that only a fraction $a_{\rm sph} \simeq 1/3$ of the B–L asymmetry ends up into a baryon asymmetry through sphaleron conversion and the dilution of the asymmetry $f \simeq 35$ due to the photon production from the time of leptogenesis till recombination. Each κ_i is the efficiency factor associated with the asymmetry production from the decays of N_i. In general the baryon asymmetry will thus depend on 10 unknown seesaw parameters: the absolute neutrino mass scale m_1, the three M_i and the 6 parameters in Ω. However, assuming a mild heavy neutrino mass hierarchy, $M_2 \gtrsim 5\, M_1$, and assuming that the inverse decays of the lightest RH neutrino N_1 strongly wash-out the asymmetry generated from the two heavier ones, then one has a simplified picture where the final asymmetry is produced only from N_1 decays and $\eta_B \simeq 10^{-2}\, \varepsilon_1\, \kappa_1$. For values of $M_1 \ll 10^{14}\,{\rm GeV}\, m_{\rm atm}^2/\sum_i m_i^2$, the main contribution to the wash-out comes from inverse decays [6, 7] and κ_1 is a function just of the *effective neutrino mass* $\widetilde{m}_1$ defined as $\widetilde{m}_1 \equiv (m_{\rm D}^\dagger\, m_{\rm D})_{11}/M_1$. The assumption of strong wash-out holds for $\widetilde{m}_1 \gtrsim m_\star \simeq 10^{-3}\,{\rm eV}$, where $m_\star$ is the *equilibrium neutrino mass*. The effective neutrino mass can be conveniently expressed in terms of the Ω matrix as $\widetilde{m}_1 = \sum_i m_i\, |\Omega_{j1}^2|$ and from the Ω orthogonality it easily follows that $\widetilde{m}_1 \geq m_1$ [8]. This is the only model independent information on $\widetilde{m}_1$, quite relevant if $m_1 \gtrsim m_\star$, since in this case one can conclude that the strong wash-out condition is always realized. If $m_1 \lesssim m_\star$, then $\widetilde{m}_1$ can also lie in the weak-wash out regime. However, now, for fully normal (inverted) hierarchical neutrinos, one has

$$\widetilde{m}_1 \simeq m_1\, |\Omega_{11}^2| + m_{\rm sol}\, (m_{\rm atm})\, |\Omega_{21}^2| + m_{\rm atm}\, |\Omega_{31}^2| \, , \tag{10}$$

and one can easily understand that the condition $\widetilde{m}_1 \ll m_\star \ll m_{\rm sol}$ holds only when [9]

$$\Omega \simeq \begin{pmatrix} 1 & 0 & 0 \\ 0 & \Omega_{22} & \sqrt{1-\Omega_{22}^2} \\ 0 & -\sqrt{1-\Omega_{22}^2} & \Omega_{22} \end{pmatrix}, \tag{11}$$

such that $|\Omega_{21}^2| \ll m_1/m_2$ and $|\Omega_{31}^2| \ll m_1/m_{\rm atm}$.

If $M_1 \gtrsim 10^{14}\,{\rm GeV}\, m_{\rm atm}^2/\sum_i m_i^2$, then off-shell $\Delta L = 2$ processes gives a non negligible contribution to the wash-out such that

$$\kappa_1 \simeq \kappa_1(\widetilde{m}_1)\, e^{-\frac{M_1}{10^{14}\,{\rm GeV}}\,\frac{\sum_i m_i^2}{m_{\rm atm}^2}}. \tag{12}$$

Another nice consequence of assuming a hierarchical heavy neutrino spectrum is that it allows to write down an approximate expression for ε_1,

$$\varepsilon_1 \simeq \varepsilon_1(M_1, m_1, \widetilde{m}_1, \Omega_{j1}^2) \equiv \overline{\varepsilon}(M_1)\ \beta(m_1, \widetilde{m}_1, \Omega_{j1}^2), \tag{13}$$

where $\overline{\varepsilon}(M_1) \equiv 3\, M_1\, m_{\rm atm}/(16\,\pi\,{\rm v}^2)$. Therefore, we can see that within the simplified picture where the dominant contribution to the final asymmetry arises from N_1 decays, one has that $\eta_B = \eta_B(M_1, m_1, \widetilde{m}_1, \Omega_{j1}^2)$ depends just on 6 parameters. One can find interesting constraints imposing that the predicted asymmetry explains the observed value. Interestingly, the function $\varepsilon_1(M_1, m_1, \widetilde{m}_1, \Omega_{j1}^2)$ has an upper bound that can be found maximizing β over the Ω_{j1}^2's [10, 11]. This is saturated for fully hierarchical neutrinos ($m_1 = 0$) and when $\Omega_{21}^2 = {\rm Re}(\Omega_{31}^2) = 0$, such that $\beta = 1$ and $\varepsilon_1 = \overline{\varepsilon}(M_1)$. The upper bound on ε_1 implies a lower bound on M_1 [11, 12]

$$M_1 \gtrsim 4.2 \times 10^8\,{\rm GeV} \times [k_1(\widetilde{m}_1)]^{-1} \quad (M_1 \ll 10^{14}\,{\rm GeV}), \tag{14}$$

and consequently a lower bound on the reheating temperature [7] $T_{\rm reh} \gtrsim \frac{M_1}{5}$. For non-zero values of m_1, the upper bound on ε_1 becomes more restrictive [11, 13],

$$\varepsilon_1 \leq \overline{\varepsilon}(M_1, m_1, \widetilde{m}_1) \equiv \overline{\varepsilon}(M_1)\,\frac{m_{\rm atm}}{m_1 + m_3}\, f(m_1, \widetilde{m}_1). \tag{15}$$

The bound is still saturated for $\Omega_{21}^2 = 0$ but this time $X \equiv {\rm Re}(\Omega_{31}^2) \neq 0$ [14]. The function $f(m_1, \widetilde{m}_1) \in [0,1]$ vanishes for $\widetilde{m}_1 = m_1$ and is equal to 1 in the limit $m_1/\widetilde{m}_1 \to 0$. For generic values of m_1 it can be calculated simply finding the maximum of $Y \equiv {\rm Im}(\Omega_{31}^2)$ for fixed $\widetilde{m}_1$ and with the orthogonality implying $\Omega_{11}^2 + \Omega_{31}^2 = 1$ and then $f = (m_1 + m_3)\, Y/\widetilde{m}_1$ [9]. For $m_1 \ll m_{\rm atm}$ the function f is well approximated by [9, 13]

$$f(m_1, \widetilde{m}_1) = \frac{m_3 - m_1\sqrt{1 + \frac{m_3^2 - m_1^2}{\widetilde{m}_1^2}}}{m_3 - m_1}, \tag{16}$$

while in the limit of quasi-degenerate neutrinos one has [9, 14]

$$f = \sqrt{1 - \left(\frac{m_1}{\widetilde{m}_1}\right)^2} \,. \tag{17}$$

It is quite remarkable that, accounting also for the suppression of the efficiency factor (*cf.*(12)), the final baryon asymmetry is strongly suppressed for $m_1 \gg m_{\rm atm}$. This results in a stringent upper bound on the neutrino masses $m_i \leq 0.1\,{\rm eV}$ [13]. For a generic choice of the seesaw parameters, the CP asymmetry can be written as

$$\varepsilon_1 = \overline{\varepsilon}(M_1, m_1, \widetilde{m}_1)\, \sin\delta_{\rm L}(m_1, \widetilde{m}_1, \Omega^2_{j1})\,, \tag{18}$$

where the *effective leptogenesis phase* $\sin\delta_{\rm L}$ takes into account the suppression of the asymmetry compared to the case when the upper bound is saturated and $\sin\delta_{\rm L} = 1$. The exponential suppression of κ_1 for large M_1 (*cf.* (12)) yields a lower bound on $\sin\delta_{\rm L}$ given by [9]

$$\sin\delta_{\rm L} \gtrsim 4\times 10^{-2}\,(\widetilde{m}_1/{\rm eV}) \quad (\widetilde{m}_1 \gtrsim 10^{-3}\,{\rm eV})\,. \tag{19}$$

All the derived bounds are valid under a set of minimal assumptions and it is actually non trivial that the final asymmetry can be explained within such a minimal picture. In particular it is interesting that this is possible only because the atmospheric neutrino mass scale lies in the range $10^{-3}\,{\rm eV} \lesssim m_{\rm atm} \lesssim 1\,{\rm eV}$ [15], a result to be regarded as a successful test for the minimal leptogenesis scenario.

On the other hand the existence of the constraints on neutrino masses, the lower bound on the reheating temperature of the Universe and the upper bound on the neutrino masses, can potentially disprove the minimal scenario. In particular the first is a problem when leptogenesis is embedded within the minimal supersymmetric standard model scenario where an upper bound on $T_{\rm reh}$ is necessary to avoid a gravitino over-production. The second is, of course, a problem if neutrino masses larger than $0.1\,{\rm eV}$ will be found. At the moment it is intriguing that the most restrictive upper bound from a combined analysis of the CMB acoustic peaks and of large scale structure gives $m_1 < 0.14\,{\rm eV}$ and thus it is in agreement with the bound from leptogenesis. Therefore, at the moment, the minimal picture cannot be ruled out.

3. A new scenario of leptogenesis

In the weak wash-out regime, for $\widetilde{m}_1 \lesssim m_\star$, the minimal picture where the asymmetry is produced by the N_1's, encounters two serious obstacles. The first is that the CP asymmetry ε_1 vanishes in the limit $\widetilde{m}_1 \to m_1$, implying that one has to tune $\widetilde{m}_1$ such that it is lower than $m_\star \simeq 10^{-3}\,\mathrm{eV}$ but much larger than m_1. The second problem is that, in the weak wash-out regime, the final asymmetry depends on the initial asymmetry and, more importantly, on the initial N_1 abundance and a calculation of the final asymmetry becomes very model dependent [7].

These two problems are both solved considering another possibility: that the asymmetry is produced from the decays of N_2, the next-to-lightest RH neutrino. Indeed if $\widetilde{m}_1 \lesssim m_\star$, as we are assuming, then necessarily $\widetilde{m}_2 \gtrsim m_{\rm sol} \gg m_\star$ and thus the generation occurs in the strong wash-out regime, solving the problem of the initial conditions. Moreover now if $\widetilde{m}_1 = m_1$, when Ω is given by the Eq. (11), the CP asymmetry ε_2 does not vanish. Therefore one does not need to fine tune Ω, it is enough that this is close to the form given by Eq. (11) and, of course, that the phase of Ω_{22} is large enough.

Another attractive feature is that N_1 does not play any role and thus the lower bound on M_1 disappears and is replaced by an analogous lower bound on M_2, still implying a lower bound on $T_{\rm reh} \gtrsim M_2/5$. It is intriguing to speculate about *possible phenomenological implications of a light* N_1. The first is to think whether this can be light enough to be produced in the Large Hadron Collider. However, the N_1 Yukawa coupling upper bound, $U^\dagger h_{11} \lesssim 10^{-7}\,\sqrt{M_1/1\,\mathrm{TeV}}$, gives no hopes, unless new extra-gauge interactions are assumed.

4. Beyond the minimal picture

The new scenario described in the previous section is an interesting logical completion of the minimal picture. However, the possibility to evade the lower bound on the reheating temperature and the upper bound on the neutrino masses requires more drastic departures from the minimal scenario. A large variety of models have been proposed. They can be classified in three categories.

Leptogenesis from type II seesaw formula [16]. A first class of models is based on modifications of the minimal seesaw formula. The most popular is the account of the triplet Higgs term in the neutrino mass term (*cf.* (2)). In this, the simple seesaw formula (*cf.* (3) gets generalized into

$$m_\nu = -m_{\rm D}\,\frac{1}{M}\,m_{\rm D}^T - h^T\,\frac{v^2}{M_T}\,. \tag{20}$$

The dominant wash-out processes constraints mainly the first term but not the second and in this way the upper bound on the neutrino masses can be easily evaded. Moreover, the CP asymmetry receives additional contributions making possible also to evade the lower bounds on M_1 and on $T_{\rm reh}$.

Degenerate heavy neutrino spectrum [9,13,14,17]. If $M_2 \lesssim 5\,M_1$ one expects deviations from the constraints obtained in the minimal picture, even in the strong wash-out regime for $\widetilde{m}_1 \gg m_\star$. There are two effects to consider: one is that the asymmetry and the wash-out from the $N_{2,3}$ decays and inverse decays can give a non negligible contribution and so one has to consider the general expression (9) for the final asymmetry. The second effect is that there can be significant deviations from the approximate expression for ε_1 (*cf.* (13)). Defining $\xi_\varepsilon \equiv \varepsilon_1/\overline{\varepsilon}(M_1)$, one has that for $M_2 \lesssim 5\,M_1$ values $\xi_\varepsilon \gtrsim 1.1$ are possible, with a consequent relaxation of the neutrino mass constraints. However, the possibility of a significant evasion of the bounds, especially the upper bound on the neutrino masses, is possible if the CP asymmetry undergoes a resonant enhancement, such that $\varepsilon_1 \sim 0.1$ independently on the values of M_1 and m_1 [14,17].

Non thermal leptogenesis [18]. If the temperature $T_{\rm reh} \lesssim M_1/5$, a thermal production of the N_1 is inefficient. Different mechanisms of non thermal production have been proposed, all associated to the occurrence of an inflationary stage. The bounds of the minimal picture can be easily evaded.

Within all these three categories of non minimal models, neutrino mass bounds can be easily evaded. However, at the same time the correlation between the neutrino mixing mass scales and the baryon asymmetry, the *leptogenesis conspiracy*, gets lost or in other words is like ‘to throw away the baby with the bath water’. Therefore, until data will not disprove the minimal picture, this appears the most appealing way to realize leptogenesis. In the near future the most important experimental tests for the minimal picture are represented by the measurement of the absolute neutrino mass scale to be compared with the upper bound $m_i < 0.1\,{\rm eV}$ and a test of the supersymmetric models in LHC, that could confirm the gravitino upper bound on $T_{\rm reh}$ at variance with the minimal leptogenesis lower bound.

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