PRAMANA — journal of physics | © Indian Academy of Sciences | Vol. 67, No. 4 October 2006 pp. 723–733

Leptogenesis models and neutrino mass constraints

T HAMBYE
Instituto de Física Teórica, Universidad Autónoma de Madrid, Cantoblanco, Spain
E-mail: thomas.hambye@uam.es

Abstract. Through leptogenesis, baryogenesis could have the same origin as neutrino masses. We review the various ways of implementing the leptogenesis mechanism. Emphasis is put on the conditions which, in order that this mechanism works, need to be fulfilled by the neutrino masses as well as by the heavy state masses.

Keywords. Baryogenesis; neutrino masses.

PACS Nos 14.40.Pq; 98.80.Cq

1. Introduction

Following the recent convincing evidence for neutrino masses, leptogenesis [1] has become a well motivated possible explanation of the origin of the baryon asymmetry of the universe. In these proceedings, in the framework of the usual seesaw model with three right-handed neutrinos ('Type-I' seesaw model [2]), we will study the neutrino mass constraints, in particular the neutrino mass upper bound, which exist in order that leptogenesis can successfully explain the baryon asymmetry of the universe. We will show how this upper bound, which is quite stringent for a very hierarchical spectrum of right-handed neutrino masses, can be largely relaxed by no longer assuming such a spectrum. We will also review other attractive models of generation of both neutrino masses and leptogenesis, putting special emphasis on models where leptogenesis is induced by the decay of a heavy scalar $SU(2)_L$ triplet, for which a precise calculation of the efficiency and constraints has been recently performed.

2. The three basic ingredients of leptogenesis

To introduce leptogenesis we consider the usual 'Type-I' seesaw model with three heavy right-handed neutrinos N_i. It is based on the following Lagrangian:

$$L \ni \left(\lambda^{ij} H^\dagger N_i L_j + \frac{M_{N_i}}{2} N_i N_i + \text{h.c.}\right), \tag{1}$$

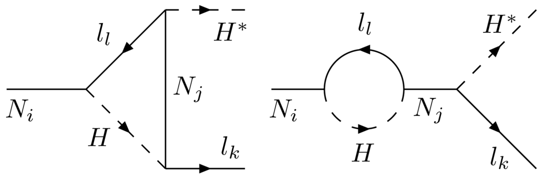

Figure 1. One-loop diagrams contributing to the asymmetry from the N_i decay.

with $L_j = (\nu_{jL}, l_{jL})^T$, $H = (H^0, H^-)^T$ (where we made the choice to work in the basis where the N_i mass matrix is real and diagonal). This model has 18 parameters: 9 combinations of them enter in the neutrino mass matrix $M_\nu^I = -\lambda^T M_N^{-1} \lambda v^2$, and nine of them decouple from it. For the following we order the neutrino masses as: $m_{\nu_3} > m_{\nu_2} > m_{\nu_1} \geq 0$ and $M_{N_3} \geq M_{N_2} \geq M_{N_1} \geq 0$.

The crucial ingredient for leptogenesis is *the CP-asymmetry*, that is to say the averaged ΔL produced each time one N_i decays (at temperature of order their mass when they disappear from the universe thermal bath by decaying). At lowest order, that is to say at one-loop order, for N_1 (and similarly for $N_{2,3}$), it is given by (see e.g. [3])

$$\varepsilon_{N_1} \equiv \frac{\Gamma(N_1 \to l H^*) - \Gamma(N_1 \to \bar{l} H)}{\Gamma(N_1 \to l H^*) + \Gamma(N_1 \to \bar{l} H)} = -\sum_{j=2,3} \frac{3}{2} \frac{M_{N_1}}{M_{N_j}} \frac{\Gamma_{N_j}}{M_{N_j}} I_j \frac{2S_j + V_j}{3} \tag{2}$$

where $I_j = \mathrm{Im}\,[(\lambda\lambda^\dagger)^2_{1j}]/|\lambda\lambda^\dagger|_{11}|\lambda\lambda^\dagger|_{jj}$ and

$$\frac{\Gamma_{N_j}}{M_{N_j}} = \frac{|\lambda\lambda^\dagger|_{jj}}{8\pi} \equiv \frac{\tilde{m}_j M_{N_j}}{8\pi v^2} \tag{3}$$

with

$$S_j = \frac{M_{N_j}^2 \Delta M_{1j}^2}{(\Delta M_{1j}^2)^2 + M_{N_1}^2 \Gamma_{N_j}^2},$$
$$V_j = 2 \frac{M_{N_j}^2}{M_{N_1}^2} \left[\left(1 + \frac{M_{N_j}^2}{M_{N_1}^2}\right) \log\left(1 + \frac{M_{N_1}^2}{M_{N_j}^2}\right) - 1 \right],$$

where $\Delta M_{ij}^2 = M_{N_j}^2 - M_{N_i}^2$. The factors $S_j(V_j)$ comes from the one-loop self-energy (vertex) contribution to the decay widths (figure 1). The I_j factors are the CP-violating coupling combinations entering in the asymmetry.

Once the averaged ΔL produced per decay has been calculated, the second ingredient to consider is *the efficiency factor* η. This factor allows to calculate the lepton asymmetry produced from the CP-asymmetry,

$$\frac{n_{\mathrm{L}}}{s} = \varepsilon_{N_i} Y_{N_i}|_{T \gg M_{N_i}} \eta\,, \tag{4}$$

where $Y_{N_i} = n_{N_i}/s$ is the number density of N_i over the entropy density, with $Y_{N_i}|_{T\gg M_{N_i}} = 135\zeta(3)/(4\pi^4 g_*)$ where $g_* = 108.5$ is the number of degrees of freedom in thermal equilibrium in the 'Type-I' model before the N_1 is decayed. If all right-handed neutrinos decay out-of-equilibrium, $\eta = 1$, but η can be much smaller than one, if they are not fully out-of-equilibrium while decaying, and/or if there are at this epoch L-violating processes partly in thermal equilibrium. The processes which can put the N_i in thermal equilibrium and/or violate L are the inverse decay process and $\Delta L = 1, 2$ scatterings. To avoid a large damping effect, it is necessary that these processes are not too fast with respect to the Hubble constant. For the inverse decay process (which is the most dangerous process, see e.g. the discussion of ref. [4]), this gives the condition: $\Gamma_{N_i}/H(T \simeq M_{N_i}) \leq 1$ with $H(T) = \sqrt{4\pi^3 g_*/45}\, T^2/M_{\rm Planck}$. In practice, to calculate η we need to put all these processes in the Boltzmann equations [5,6] which allow a precise calculation of the produced lepton asymmetry as a function of the temperature T. The corresponding efficiency factor including finite temperature effects can be found in ref. [6] in the limit where the right-handed neutrinos have a hierarchical spectrum $M_{N_1} \ll M_{N_{2,3}}$. In this limit only the asymmetry produced by the decay of the lightest right-handed neutrino N_1 survives and is important, which simplifies greatly the calculations.

Once the produced L asymmetry n_L/s has been determined, the third ingredient for leptogenesis is the calculation of the baryon asymmetry produced, due to the partial *conversion of the L asymmetry to a B asymmetry* by the Standard Model non-perturbative sphaleron processes. The conversion factor is $n_B/s = -(28/79)n_L/s$ which leads to $n_B/s = -1.38 \cdot 10^{-3}\varepsilon_{N_1}\eta$, to be compared with the experimental BBN and WMAP value: $n_B/s = (8.7 \pm 0.4) \cdot 10^{-11}$.

3. The neutrino mass constraints

In addition to the three ingredients above, there is a fourth crucial ingredient at the origin of the relevance of the leptogenesis mechanism: the neutrino mass constraints. There are two types of neutrino mass constraints on leptogenesis.

3.1 *The neutrino mass constraint on η*

The first constraint is the neutrino mass constraint on the size of the washout, i.e. on η. It comes from the fact that, in full generality, the ratio Γ_{N_i}/H, which has to be smaller than unity to have no washout suppression, is always larger than the ratio of the lightest neutrino mass m_{ν_1} over the $m_* = 16\pi^2 v^2 \sqrt{g_*\pi/45}/M_{\rm Planck} \sim 10^{-3}$ eV scale [7]:

$$\frac{\Gamma_{N_i}}{H} \geq \frac{m_{\nu_1}}{10^{-3}\,\mathrm{eV}}\,. \tag{5}$$

This inequality comes simply from eq. (3) and the $\tilde{m}_j \equiv v^2|\lambda\lambda^\dagger|^2_{jj}/M_{N_j} \geq m_{\nu_1}$ inequality. It means that, if $m_{\nu_1} > 10^{-3}$ eV, there will be some washout and the larger is m_{ν_1} above this scale the larger is the washout, i.e. the smaller is η. The fact that m_*, which is a function of the electroweak and Planck scales, is of the order of the neutrino masses. It is a quite remarkable fact that m_*, which is a function of the electroweak and Planck scales, is of the order of the neutrino masses. It means that the N_i are naturally not in deep thermal equilibrium when they decay.

3.2 *The neutrino mass constraint on* ε_{N_i}

The second neutrino mass constraint is on the size of the CP-asymmetries ε_{N_i}. It generally applies but not always, depending on both right-handed and left-handed neutrino mass spectra.

3.2.1 *Hierarchical* N_i *with hierarchical or inverted hierarchical* ν_i

If right-handed neutrino masses differ by several orders of magnitude, the size of the ε_{N_1} asymmetry is quite constrained by the size of the neutrino masses. There exists an upper bound [8–11] which in its exact form was given by ref. [11]:

$$|\varepsilon_{N_1}| \leq \frac{3}{16\pi}\frac{M_{N_1}}{v^2}\frac{\Delta m^2_{\text{atm}}}{m_{\nu_3}+m_{\nu_1}}\,. \tag{6}$$

For hierarchical or inverted hierarchical light neutrinos, $m_{\nu_3}+m_{\nu_1}$ is fixed to the value $\simeq (\Delta m^2_{\text{atm}})^{1/2}$. Due to the fact that the upper bound on ε_{N_1} is proportional to M_{N_1}, successful leptogenesis with hierarchical N_i implies a lower bound on this mass [11,8,9,6]:

$$M_{N_1} > 5\times 10^8\,\text{GeV}\,. \tag{7}$$

This result holds for the case where the N_1 are in thermal equilibrium before decaying. Starting instead (due to inflation dynamics) from a universe with no (with only) right-handed neutrinos at a temperature above their mass, this bound becomes [6]: $M_{N_1} > 2\times 10^9$ GeV (2×10^7 GeV).

Note that if the N_i masses do not differ by several orders of magnitudes but have a hierarchy similar to the ones of the charged leptons or quarks, that is to say if $M_{N_1} \simeq (10\text{–}100) M_{N_2}$ with $M_{N_3} > M_{N_2}$, the bound above can be relaxed. In this case the L-asymmetry production is still dominated by the decays of the lightest right-handed neutrino N_1 and the upper bound on the CP-asymmetry is the same as for very-hierarchical neutrinos, except that there are extra corrections [3] in $M^2_{N_1}/M^2_{N_{2,3}}$ to be added in eq. (6):

$$\delta\varepsilon_{N_1} \simeq \frac{3}{16\pi}\frac{M_{N_1}}{v^2}\tilde{m}_{2,3}\frac{M^2_{N_1}}{M^2_{N_{2,3}}}\,. \tag{8}$$

These corrections can be very large for special configurations of the Yukawa couplings. This may lead to successful leptogenesis with M_{N_1} amply smaller than 10^8 GeV [3] (see the explicit example of ref. [12] with $M_{N_1} = 10^6$ GeV).

3.2.2 *Hierarchical N_i with quasi-degenerate ν_i*

If the light neutrino masses are quasi-degenerate, i.e. with masses larger than $\simeq (\Delta m^2_{\rm atm})^{1/2}$, and if $M_{N_1} \ll M_{N_{2,3}}$, the bound of eq. (6) still holds. Since $\Delta m^2_{\rm atm} \simeq 2 \cdot 10^{-3}$ eV^2 is fixed experimentally, this bound decreases as the neutrino masses increase and get more degenerate. Therefore as the neutrino masses increase there are two suppression effects arising: the washout effect increases (eq. (5)), and the upper bound on the asymmetry decreases. Successful leptogenesis implies therefore a lower bound on M_{N_1} larger than in eq. (7) and leads to the upper bound [13,6,3]:

$$m_{\nu_3} < 0.12 - 0.15 \text{ eV}. \tag{9}$$

Note that in order to derive this bound, eq. (6) is not sufficient. As the final produced lepton asymmetry depends not only on m_{ν_3} and M_{N_1} but also, through the efficiency factor, on Γ_{N_1} (or $\tilde{m}_1$), it is necessary [13,3] to have an upper bound for fixed values of these three parameters and not only as a function of m_{ν_3} and M_{N_1} as in eq. (6). This bound can be found in ref. [3]. Note also that here too, for moderate hierarchies of right-handed neutrino mass, the corrections of eq. (8) (which unlike in eq. (6), do not necessarily decrease when m_{ν_3} increases) can be large and leads to less stringent bounds.

An important comment on the case with hierarchical N_i but quasi-degenerate ν_i is that it is not a natural case and therefore it is of limited interest. From the seesaw formula it is very difficult to get a quasi-degenerate spectrum of light neutrinos from a hierarchical N_i spectrum. This would require precise correlations (i.e. cancellations) between the Yukawa couplings and the right-handed neutrino masses (in order that the ratio between the M_{N_i} is compensated in the seesaw mass formula by the same ratio between the Yukawa couplings). If light neutrinos are quasi-degenerate, much more natural would be that the N_i are quasi-degenerate too. Therefore, for quasi-degenerate light neutrinos, the relevant bounds are not the ones of eqs (7) and (9) but the much less stringent ones of the next section.

3.2.3 *Quasi-degenerate N_i*

If at least two right-handed neutrinos have masses very close to each other, $M_{N_1} \sim M_{N_2}$, the situation changes drastically with respect to the two previous cases. This is due to the fact that the one loop self-energy diagram displays in this case a resonance behaviour [14,15,3,16], coming from the propagator of the virtual right-handed neutrino in this diagram. This effect can be seen from the S_j factors in eq. (2). Since in the seesaw model the decay widths of the N_i are generically much smaller than their masses, this resonance effect can lead to several order of magnitude enhancement of the asymmetry. At the resonance, that is to say for $M_{N_2} - M_{N_1} = \Gamma_{N_2}/2$, the S_2 factor, which is unity in the $M_{N_1} \ll M_{N_{2,3}}$ limit, is as large as $\frac{1}{2} M_{N_2}/\Gamma_{N_2}$. In this case one can show that the asymmetry is not bounded anymore by an expression depending on the neutrino masses, i.e. not proportional anymore to $\Delta m^2_{\rm atm}/(m_{\nu_3} + m_{\nu_1})$ as in eq. (6), but just by $\varepsilon_{N_1} \leq \frac{1}{2}$

[3]. Together with ε_{N_2}, which is equal to ε_{N_1} in this case and has also to be taken into account, CP-violation is just bounded by unity. As a result, since neither the maximal asymmetry nor the washout effect (coming from inverse decay) depend on M_{N_1}, successful leptogenesis can be obtained at any scale except that the L to B conversion from sphalerons still needs to be effective. This requires M_{N_1} to be above the electroweak scale (i.e. typically above ~ 1 TeV [15]). Moreover since the upper bound on the CP asymmetry is independent of neutrino masses, there is no more suppression of the asymmetry for large neutrino masses. The only remaining suppression effect arising for large neutrino masses is the one in §3.1 above coming from the washout. Therefore the upper bound on the neutrino masses in this case gets considerably relaxed. This is shown in figure 6 of ref. [3] where is plotted, as a function of m_{ν_3}, the level of degeneracy which is needed to have successful leptogenesis. Values far above the eV are possible which means that *in full generality there is no more relevant upper bound on neutrino masses coming from leptogenesis.* The value $m_{\nu_3} \simeq 1$ eV can lead to successful leptogenesis with a level of degeneracy of order $(M_{N_2} - M_{N_1})/M_{N_2} \simeq 4 \cdot 10^{-2}$ which is quite moderate (and much smaller than the level of degeneracy of the light-neutrinos for this value of m_{ν_3}). For quasi-degenerate light neutrinos, to assume specific flavor pattern, such as a weakly broken $SO(3)$ symmetry to explain the quasi-degeneracy of both heavy and light neutrinos, may lead to more stringent bounds but even in this case, a value of m_{ν_3} as large as ~ 1 eV appear to be perfectly compatible with successful leptogenesis [3].

4. Other leptogenesis models

In order to explain the neutrino masses, the Type-I seesaw mechanism is probably the most direct extension of the Standard Model and in this sense it is the most attractive. However it is not the only attractive seesaw model. Beside the Type-I seesaw, one can think about two other basic seesaw mechanisms to induce neutrino masses. The first one is the Type-II seesaw model [2] where neutrino masses are due to the exchange of a heavy scalar Higgs triplet. The second one is from the exchange of three heavy self-conjugated $SU(2)_L$ triplets of fermions [17]. In addition to these three seesaw basic mechanisms one can also think about combinations of them, such as a model with both three N_i's and one scalar triplet Δ_L, naturally present in left–right or $SO(10)$ models. Models with two scalar triplets or extra-singlets are also possible. All these models can induce baryogenesis successfully via leptogenesis, except the Type-II model with just a single scalar triplet [18,3,19], though its SUSY version can be successful from additional soft terms [19,20]. The triplet of fermion model [3] induces leptogenesis in a way very similar to the Type-I model, with however slightly more stringent bounds on the masses of the light and heavy states (due to the fact that these fermions have additional gauge scatterings). The Type-I plus Type-II models [21–24], multiple Type-II [18] or multiple Type-I models [25–28] involve new types of diagrams, which lead to no more bounds on neutrino masses [23,28], even for a hierarchical spectrum of heavy states. Note nevertheless that in all these models successful leptogenesis still leads to a lower bound on the mass of the lightest heavy states (between 10^7 GeV and 10^{11} GeV depending on the model [23,24,28]). Note also that in the supersymmetric Type-I models, soft terms can lead naturally to a resonance of the CP-asymmetry [29], which allows M_{N_1} to be orders of magnitude below 10^8 GeV. Recently it has been shown [30] that successful leptogenesis can be achieved too from the decay of right-handed neutrinos to charged right-handed leptons if a $SU(2)_L$ singlet charged scalar particle δ^+ exists. This scenario might occur in the framework of various GUT models and can induce leptogenesis at a scale as low as few TeV without resonance. Finally note that radiative models of neutrino masses at TeV scale cannot lead in general to successful leptogenesis; note nevertheless the two exceptions of refs [16,31] in the framework of the seesaw extended MSSM with neutrino masses from soft terms. See also another exception based on three-body decays in the context of the Zee model [9]. In the following we will give some details only for the Type-I plus Type-II model.

5. The Type-I plus Type-II model

The case where we add to the Standard Model three right-handed neutrinos and one scalar triplet is quite interesting because it is the situation of the ordinary left–right models and of the renormalizable $SO(10)$ models such as defined in ref. [32] (i.e. with a **126** multiplet in order to give mass to the N_i's in a renormalizable way). In this case the relevant Lagrangian is just the sum of the Lagrangian of eq. (1) and of the following triplet terms:

$$-M_\Delta^2 \mathrm{Tr}\Delta_L^\dagger \Delta_L - (Y_\Delta)_{ij} L_i^T C i\tau_2 \Delta_L L_j + \mu H^T i\tau_2 \Delta_L H + \text{h.c.}, \tag{10}$$

with Δ_L the scalar triplet in its matrix form ($\Delta_{L_{11}} = \Delta_{L_{22}} = \delta^+/\sqrt{2}$, $\Delta_{L_{12}} = \delta^{++}$, $\Delta_{L_{21}} = \delta^0$). Equation (10) leads to the neutrino mass matrix $M_\nu^{\mathrm{II}} = 2Y_\Delta v_\Delta \simeq 2Y_\Delta \mu^* v^2/M_\Delta^2$, but it alone cannot lead to successful leptogenesis (see e.g. [18,3,19]). However the simultaneous presence of both Type-I and Type-II sources of neutrino masses leads to new interesting leptogenesis contributions. To discuss this model it is necessary to consider two cases, depending on which particle is the lightest one, the scalar triplet Δ_L or the lightest right-handed neutrino N_1.

5.1 *The* $M_\Delta < M_{N_1}$ *case*

If the triplet is lighter than N_1, the production of the asymmetry will be naturally dominated by the decay of the triplet to two leptons. The CP-asymmetry comes in this case from the difference between the decay width of Δ_L^* to two leptons and of Δ_L to two anti-leptons. The leptogenesis one loop diagram is a vertex diagram involving both the decaying triplet and a virtual right-handed neutrino (see first graph of figure 2). This diagram was first displayed in ref. [21]. Calculating explicitly its contribution we get [23,33]:

$$\varepsilon_\Delta = 2 \cdot \frac{\Gamma(\Delta_L^* \to l + l) - \Gamma(\Delta_L \to \bar{l} + \bar{l})}{\Gamma_{\Delta_L^*} + \Gamma_{\Delta_L}} \tag{11}$$

$$= -\frac{1}{4\pi}\frac{M_\Delta^2}{(\sum_{ij}|(Y_\Delta)_{ij}|^2 M_\Delta^2 + |\mu|^2)}\frac{1}{v^2}\mathrm{Im}[(M_\nu^{\mathrm{I}*})_{il}(Y_\Delta)_{il}\mu^*]\,, \tag{12}$$

where M_ν^{I} is the Type-I contribution to the neutrino mass matrix (given in §2). For each of the three triplet components the total decay width is

$$\Gamma_\Delta = \frac{1}{8\pi}M_\Delta\Big(\sum_{ij}|(Y_\Delta)_{ij}|^2 + \frac{|\mu|^2}{M_\Delta^2}\Big)\,. \tag{13}$$

From this CP-asymmetry leptogenesis works in a way similar to the Type-I model apart from three important differences. The first is that, unlike the N_i which are singlets, for the triplets there are gauge scatterings which tend to put the Δ_L closer to thermal equilibrium. The second difference is that for the decay of a single triplet, there is no one-loop self-energy diagram, so there is no possible resonance effect. The third one is that the interplay between the neutrino masses and the size of the washout and the size of the asymmetry is completely different from the Type-I case. The decay width is proportional to the triplet couplings but the asymmetry is essentially proportional to the right-handed neutrino couplings (taking into account the fact that there are two triplet couplings in both numerator and denominator of the asymmetry). As a result one may consider the possibility that the Type-II contribution to neutrino masses is small enough to avoid large washout and that the Type-I contribution to neutrino masses is large. By increasing in this way the Type-I neutrino mass contribution, the washout remains unchanged but the asymmetry increases proportionally to the neutrino masses [23]. Therefore, there is no upper bound on neutrino masses in this model to have successful leptogenesis [23]. Note that there is nevertheless a lower bound on the triplet mass because the asymmetry is proportional to it. Recently a full calculation of the efficiency, including the effects of the gauge scatterings, has been performed in ref. [33]. The allowed ranges of parameters have been determined. The resulting bounds on M_Δ depends on both Type-I and Type-II contributions to the neutrino masses. For a hierarchical spectrum of light neutrinos dominated by the Type-I contribution, with Type-II contribution below 10^{-3} eV, the bound is

$$M_\Delta > 2.8\cdot 10^{10}\,\mathrm{GeV} \tag{14}$$

while for larger Type-II contribution the M_Δ lower bound increases: for a Type-II contribution responsible for the atmospheric splitting the bound is $M_\Delta > 1.3\cdot 10^{11}$ GeV. As there is no possible resonance effect these bounds are absolute bounds for hierarchical light neutrinos. For larger values of m_{ν_3} these bounds decrease because as explained above in this case the asymmetry increases. However, assuming m_{ν_3} below 1 eV, this bound cannot decrease by more than one order of magnitude.

5.2 *The $M_\Delta > M_{N_1}$ case*

If the triplet is sizeably heavier than at least one right-handed neutrino, then it is the decay of the lightest right-handed neutrino to a lepton and a Higgs boson which

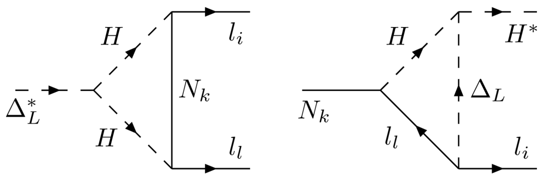

Figure 2. Diagrams contributing to the asymmetry from the Δ_L and N decays.

dominates the production of the L asymmetry. In this case leptogenesis can be produced from the pure Type-I model just as in §§2 and 3. However, in addition to this pure Type-I contribution, there is a new contribution to leptogenesis coming from a diagram [21–24] with a real right-handed neutrino and a virtual triplet (second diagram of figure 2). Disregarding the pure Type-I contribution (assuming that it has a small contribution because the Type-I Yukawa couplings and/or their phases are small) this new diagram can perfectly lead to successful leptogenesis alone. For the asymmetry defined in eq. (2) one obtains from this diagram [23,24]

$$\varepsilon^{\Delta}_{N_1} = \frac{3}{16\pi}\frac{M_{N_1}}{v^2}\frac{\sum_{il}\mathrm{Im}[\lambda_{1i}\lambda_{1l}(M_\nu^{\mathrm{II}*})_{il}]}{\sum_i|\lambda_{1i}|^2}\,, \tag{15}$$

where M_ν^{II} is the neutrino mass matrix induced by the Type-II contribution given above. The discussion is similar to the one in the $M_\Delta < M_{N_1}$ case inverting the role of Type-I and Type-II two contributions. It is now the decay of the right-handed neutrino to lepton and Higgs, induced by the Type-I couplings (eq. (3)), which essentially determines the washout. To increase the asymmetry without increasing the washout, one can then consider the possibility of keeping the Type-I contribution small, increasing the Type-II contribution which increases the asymmetry but not the washout. As a result, here too, there is no more upper bound on neutrino masses for leptogenesis [23]. For hierarchical light neutrinos the lower bound on the lightest right-handed neutrino mass is to a good approximation the same as in the pure Type-I model [23] (eq. (7)). As the asymmetry is linear in both M_{N_1} and the neutrino masses, for larger values of m_{ν_3} the asymmetry linearly increases, so the bound on M_{N_1} linearly decreases. But here too assuming m_{ν_3} below 1 eV, this bound cannot decrease by more than one order of magnitude. The precise bound is given in ref. [24] as a function of the efficiency factor η.

Note that, if instead of considering a dominant Type-II contribution, we consider the case where both Type-I and Type-II contributions are important, this lower bound on M_{N_1} does not get relaxed. In this case both contributions appear to be just proportional to their respective contributions to the neutrino masses so that, barring a possible cancellation of CP-violating phases, leptogenesis is expected to be dominated by the contribution which dominates the neutrino masses [23]. From eq. (2), the pure Type-I contribution to leptogenesis for hierarchical right-handed neutrinos turns out to be given by eq. (15) replacing M_ν^{II} by M_ν^{I}, a fact which can nicely be understood by using effective dimension five neutrino mass operators [24].

For a quasi-degenerate spectrum of right-handed neutrinos the pure Type-I contribution to leptogenesis is expected to be dominant and the triplet contribution which does not display any resonant behavior is in this case negligible.

6. Summary

In summary, there are quite a few models which from the same interactions can lead to successful generation of neutrino masses and leptogenesis. The leptogenesis upper bound on neutrino masses is quite sensitive to the model considered as well as on the assumptions made on the heavy mass spectrum in each model. In the Type-I model (and similarly in the triplet of fermion model) a stringent bound (eq. (9)), can be found only assuming a hierarchical spectrum of right-handed neutrinos. However this assumption does not appear to be the most natural assumption one can make when considering this bound, for which the light neutrinos have a quasi-degenerate spectrum. A more natural quasi-degenerate spectrum of right-handed neutrinos leads instead in full generality to an upper bound on neutrino masses far beyond the eV scale. Only with extra assumptions (such as low reheating temperature or with particular flavor structures) one might get a more stringent bound. In the other models (multiple Type-I and/or Type-II) the upper bound is far above the eV scale even for a hierarchical spectrum of heavy states.

Similarly the lower bound on the scale of leptogenesis is sensitive to the model and/or the heavy mass spectrum considered. However for this scale there is at least one firm conclusion one can draw: in all the seesaw models we considered here, if the masses of the heavy states differ by orders of magnitude, this scale has to be orders of magnitude above the TeV scale, that is to say above 10^7–10^{11} GeV depending on the model.

References

[1] M Fukugita and T Yanagida, *Phys. Lett.* **B174**, 45 (1986)
[2] For a list of references for Type-I and Type-II models see ref. [28]
[3] T Hambye, Y Lin, A Notari, M Papucci and A Strumia, *Nucl. Phys.* **B695**, 169 (2004)
[4] P Di Bari, hep-ph/0406115
[5] See, for e.g., M Plumacher, *Baryon asymmetry, neutrino mixing and supersymmetric SO(10)*, Ph.D. thesis, hep-ph/9807557
[6] G F Giudice, A Notari, M Raidal, A Riotto and A Strumia, *Nucl. Phys.* **B685**, 89 (2004)
[7] W Buchmüller and M Plumacher, *Phys. Rep.* **320**, 329 (1999)
[8] R Barbieri, P Creminelli, A Strumia and N Tetradis, *Nucl. Phys.* **B575**, 61 (2000)
W Buchmüller and T Yanagida, *Phys. Lett.* **B445**, 399 (1999)
[9] T Hambye, *Nucl. Phys.* **B633**, 171 (2002)
[10] K Hamaguchi, H Murayama and T Yanagida, *Phys. Rev.* **D65**, 043512 (2002)
[11] S Davidson and A Ibarra, *Phys. Lett.* **B535**, 25 (2002)
[12] M Raidal, A Strumia and K Turzyński, hep-ph/0408015
[13] W Buchmüller, P Di Bari and M Plümacher, *Nucl. Phys.* **B665**, 445 (2003)
[14] M Flanz, E A Paschos and U Sarkar, *Phys. Lett.* **B345**, 248 (1995)
M Flanz, E A Paschos, U Sarkar and J Weiss, *Phys. Lett.* **B389**, 693 (1996)
L Covi, E Roulet and F Vissani, *Phys. Lett.* **B384**, 169 (1996)
[15] A Pilaftsis, *Phys. Rev.* **D56**, 5431 (1997)
A Pilaftsis and T E J Underwood, hep-ph/0309342
[16] T Hambye, J March-Russell and S West, *J. High Energy Phys.* **0407**, 070 (2004)

[17] R Foot, H Lew, X-G He and G C Joshi, *Z. Phys.* **C44**, 441 (1989)
E Ma, *Phys. Rev. Lett.* **81**, 1171 (1998)
[18] E Ma and U Sarkar, *Phys. Rev. Lett.* **80**, 5716 (1998)
T Hambye, E Ma and U Sarkar, *Nucl. Phys.* **B602**, 23 (2001)
[19] G D'Ambrosio, T Hambye, A Hektor, M Raidal and A Rossi, *Phys. Lett.* **B604**, 199 (2004)
[20] E J Chun and S Scopel, *Phys. Lett.* **B636**, 278 (2006)
[21] P O'Donnell and U Sarkar, *Phys. Rev.* **D49**, 2118 (1994)
[22] G Lazarides and Q Shafi, *Phys. Rev.* **D58**, 071702 (1998)
[23] T Hambye and G Senjanovic, *Phys. Lett.* **B582**, 73 (2004)
[24] S Antusch and S F King, *Phys. Lett.* **B597**, 199 (2004)
[25] C H Albright and S M Barr, *Phys. Rev.* **D70**, 033013 (2004)
[26] U Sarkar, hep-ph/0409019
[27] A Abada, H Aissaoui and M Losada, hep-ph/0409343
[28] T Hambye, hep-ph/0412053
[29] L Boubekeur, hep-ph/0208003
Y Grossman, T Kashti, Y Nir and E Roulet, *Phys. Rev. Lett.* **91**, 251801 (2003)
G D'Ambrosio, G F Giudice and M Raidal, *Phys. Lett.* **B575**, 75 (2003)
[30] M Frigerio, T Hambye and E Ma, hep-ph/0603123
[31] L Boubekeur, T Hambye and G Senjanovic, *Phys. Rev. Lett.* **93**, 111601 (2004)
[32] B Bajc, G Senjanovic and F Vissani, *Phys. Rev. Lett.* **90**, 051802 (2003)
R N Mohapatra and S P Ng, *Phys. Lett.* **B570**, 215 (2003)
[33] T Hambye, M Raidal and A Strumia, *Phys. Lett.* **B632**, 667 (2006)