NEUTRINOS AND MATTER-ANTIMATTER ASYMMETRY
OF THE UNIVERSE

WILFRIED BUCHMÜLLER

Deutsches Elektronen-Synchrotron DESY, 22603 Hamburg, Germany

ABSTRACT

Interactions of heavy Majorana neutrinos in the thermal phase of the early universe may be the origin of the cosmological matter-antimatter asymmetry. Successful baryogenesis, independent of initial conditions, is possible for neutrino masses in the range $10^{-3} \text{eV} \leq m_i \lesssim 0.1 \text{eV}$. Remarkably, this mass window is consistent with the evidence for neutrino masses from oscillations.

1. Matter-Antimatter Asymmetry

One of the main successes of the standard early-universe cosmology is the prediction of the abundances of the light elements, $\text{D, } ^3\text{He, } ^4\text{He and } ^7\text{Li}$. Agreement between theory and observation is obtained for a certain range of the parameter $\eta_B$, the ratio of baryon density and photon density$^\text{[1]}$,

$$\eta_B^{BBN} = \frac{n_B}{n_\gamma} = (2.6 - 6.2) \times 10^{-10},$$

where the present number density of photons is $n_\gamma \sim 400/\text{cm}^3$. Since no significant amount of antimatter is observed in the universe, the baryon density coincides with the cosmological baryon asymmetry, $\eta_B = (n_B - n_\bar{B})/n_\gamma$.

The precision of measurements of the baryon asymmetry has dramatically improved with the observation of the acoustic peaks in the cosmic microwave background radiation (CMB). Most recently, the WMAP Collaboration has measured the baryon asymmetry with a (1$\sigma$) standard error of $\sim 5\%$$^\text{[2]}$,

$$\eta_B^{CMB} = (6.1^{+0.3}_{-0.2}) \times 10^{-10}. $$

A matter-antimatter asymmetry can be dynamically generated in an expanding universe if the particle interactions and the cosmological evolution satisfy Sakharov’s conditions$^\text{[3]}$.

$^\text{a}$Presented at ‘Neutrino Telescopes’, Venice, March 2003
• baryon number violation ,
• $C$ and $CP$ violation ,
• deviation from thermal equilibrium .

Although the baryon asymmetry is just a single number, it provides an important relationship between the standard model of cosmology, i.e. the expanding universe with Robertson-Walker metric, and the standard model of particle physics as well as its extensions.

At present there exist a number of viable scenarios for baryogenesis\(^4\)\(^5\). They can be classified according to the different ways in which Sakharov’s conditions are realized. In grand unified theories baryon number ($B$) and lepton number ($L$) are broken by the interactions of gauge bosons and leptoquarks. This is the basis of classical GUT baryogenesis\(^6\). Analogously, the lepton number violating decays of heavy Majorana neutrinos lead to leptogenesis\(^6\). In the simplest version of leptogenesis the initial abundance of the heavy neutrinos is generated by thermal processes. Alternatively, heavy neutrinos may be produced in inflaton decays, in the reheating process after inflation, or by brane collisions. The observed magnitude of the baryon asymmetry can be obtained for realistic neutrino masses.

The crucial deviation from thermal equilibrium can also be realized in several ways. One possibility is a sufficiently strong first-order electroweak phase transition which would make electroweak baryogenesis possible. For the classical GUT baryogenesis and for thermal leptogenesis the departure from thermal equilibrium is due to the deviation of the number density of the decaying heavy particles from the equilibrium number density. How strong this departure from equilibrium is depends on the lifetime of the decaying heavy particles and the cosmological evolution.

A crucial ingredient of baryogenesis is the connection between baryon number and lepton number in the high-temperature, symmetric phase of the standard model. Due to the chiral nature of the weak interactions $B$ and $L$ are not conserved\(^7\). At zero temperature this has no observable effect due to the smallness of the weak coupling. However, as the temperature approaches the critical temperature $T_c$ of the electroweak phase transition, $B$ and $L$ violating processes come into thermal equilibrium\(^8\).

The rate of these processes is related to the free energy of sphaleron-type field configurations which carry topological charge. In the standard model they lead to an effective interaction of all left-handed fermions\(^7\) (cf. fig. \(1\)),

$$O_{B+L} = \prod_i (q_i^L q_i^L q_i^L l_i^L) ,$$

which violates baryon and lepton number by three units,

$$\Delta B = \Delta L = 3 .$$
Figure 1: One of the 12-fermion processes which are in thermal equilibrium in the high-temperature phase of the standard model.

The sphaleron transition rate in the symmetric high-temperature phase has been evaluated by combining an analytical resummation with numerical lattice techniques. The result is, in accord with previous estimates, that $B$ and $L$ violating processes are in thermal equilibrium for temperatures in the range

$$T_{EW} \sim 100 \text{ GeV} < T < T_{SPH} \sim 10^{12} \text{ GeV}.$$  \hspace{1cm} (5)

Sphaleron processes have a profound effect on the generation of the cosmological baryon asymmetry. Eq. (4) suggests that any $B + L$ asymmetry generated before the electroweak phase transition, i.e., at temperatures $T > T_{EW}$, will be washed out. However, since only left-handed fields couple to sphalerons, a non-zero value of $B + L$ can persist in the high-temperature, symmetric phase if there exists a non-vanishing $B - L$ asymmetry. An analysis of the chemical potentials of all particle species in the high-temperature phase yields the following relation between the baryon asymmetry and the corresponding $L$ and $B - L$ asymmetries,

$$\langle B \rangle_T = c_S \langle B - L \rangle_T = \frac{c_S}{c_S - 1} \langle L \rangle_T.$$  \hspace{1cm} (6)

Here $c_S$ is a number $\mathcal{O}(1)$. In the standard model with three generations and one Higgs doublet one has $c_s = 28/79$.

An important ingredient in the theory of baryogenesis is also the nature of the electroweak transition. A first-order phase transition yields a departure from thermal equilibrium. Since in the standard model baryon number, $C$ and $CP$ are not conserved, it is conceivable that the cosmological baryon asymmetry has been generated at the electroweak phase transition. Detailed studies during the past years
have shown that for Higgs masses above the present LEP bound of 114 GeV electroweak baryogenesis is not viable, except for some supersymmetric extensions of the standard model\(^\text{10}\). In particular, the electroweak transition may have been just a smooth crossover, without any departure from thermal equilibrium. In this case it’s sole effect in the cosmological evolution has been to switch off the \( B - L \) changing sphaleron processes adiabatically.

Based on the relation (6) between baryon and lepton number we then conclude that \( B - L \) violation is needed to explain the cosmological baryon asymmetry if baryogenesis took place before the electroweak transition, i.e. at temperatures \( T > T_{EW} \sim 100 \text{ GeV} \). In the standard model, as well as its supersymmetric version and its unified extensions based on the gauge group \( SU(5) \), \( B - L \) is a conserved quantity. Hence, no baryon asymmetry can be generated dynamically in these models and one has to consider extensions with lepton number violation\(^b\).

The remnant of lepton number violation at low energies is the appearance of an effective \( \Delta L = 2 \) interaction between lepton and Higgs fields (cf. fig. 2),

\[
\mathcal{L}_{\Delta L=2} = \frac{1}{2} f_{ij} l_i^T \phi C l_j \phi + \text{h.c.} .
\]

Such an interaction arises in particular from the exchange of heavy Majorana neutrinos. In the Higgs phase of the standard model, where the Higgs field acquires a vacuum expectation value, it gives rise to Majorana masses of the light neutrinos \( \nu_e \), \( \nu_\mu \) and \( \nu_\tau \).

Lepton number violation appears to be necessary to understand the cosmological baryon asymmetry. However, it can only be weak, since otherwise any baryon asymmetry would be washed out. The interplay of these conflicting conditions leads to important contraints on neutrino properties and on extensions of the standard model in general.

\(^b\)In the case of Dirac neutrino masses, where the Yukawa couplings of right-handed neutrinos are very small, one can construct models where an asymmetry of lepton doublets is accompanied by an asymmetry of right-handed neutrinos such that the total lepton number is conserved and \( \langle B - L \rangle_T = 0 \)\(^\text{11}\text{12}\).
2. Leptogenesis

Lepton number is naturally violated in grand unified theories (GUTs). The unification of gauge couplings at high energies suggests that the standard model gauge group is part of a larger simple group,

\[ G_{SM} = U(1) \times SU(2) \times SU(3) \subset SU(5) \subset SO(10) \ldots . \]  

The simplest GUT is based on the gauge group SU(5)\[13\). Here quarks and leptons are grouped into the multiplets,

\[ 10 = (q_L, u^c_R, e^c_R), \quad 5^* = (d^c_R, l_L), \quad (1 = \nu_R). \]  

Unlike gauge fields, quarks and leptons are not unified in a single multiplet. In particular, right-handed neutrinos are not needed in SU(5) models. Since they are singlets with respect to SU(5), they can have Majorana masses \( M \) which are not controlled by the Higgs mechanism.

The three SU(5) multiplets can have Yukawa interactions with two Higgs fields, \( H_1(5) \) and \( H_2(5^*) \),

\[ \mathcal{L} = h_{uij} 10_i 10_j H_1(5) + h_{dij} 5^*_i 10_j H_2(5^*) + h_{\nu ij} 5^*_i 1_j H_1(5) + M_{ij} 1_i 1_j. \]  

Electroweak symmetry breaking then leads to quark and charged lepton mass matrices and to the Dirac neutrino mass matrix \( m_D = h_\nu v_1 \), where \( v_1 = \langle H_1 \rangle \). The Majorana mass term, which violates lepton number (\( \Delta L = 2 \)), is invariant under SU(5), and the Majorana masses can therefore be much larger than the electroweak scale, \( M \gg v \).

All quarks and leptons of one generation are unified in a single multiplet in the GUT group SO(10)\[14\],

\[ 16 = 10 + 5^* + 1. \]  

Right-handed neutrinos are now required by the fundamental gauge symmetry, and the theory contains all ingredients needed to account for the recent evidence for neutrino masses and mixings. The seesaw mechanism\[15\] explains the smallness of the light neutrino masses by the largeness of the heavy Majorana masses \( M \). The theory predicts six Majorana neutrinos as mass eigenstates, three heavy (\( N \)) and three light (\( \nu \)),

\[ N \simeq \nu_R + \nu^c_R : \quad m_N \simeq M; \quad \]  

\[ \nu \simeq \nu_L + \nu^c_L : \quad m_\nu = -m_D \frac{1}{M} m_D^T. \]
In the simplest pattern of symmetry breaking, $B - L$, a subgroup of SO(10), is broken at the unification scale $L_{\text{GUT}}$. If Yukawa couplings of the third generation are $O(1)$, as it is the case for the top-quark, one finds for the corresponding heavy and light neutrino masses: $M_3 \sim L_{\text{GUT}} \sim 10^{15}$ GeV and $m_3 \sim v^2 / M_3 \sim 0.01$ eV. It is very remarkable that the light neutrino mass $m_3$ is of the same order as the mass differences $(\Delta m^2_{\text{sol}})^{1/2}$ and $(\Delta m^2_{\text{atm}})^{1/2}$ inferred from neutrino oscillations. This suggests that, via the seesaw mechanism, neutrino masses indeed probe the grand unification scale! The difference of the observed mixing patterns of quarks and leptons is a puzzle whose solution has to be provided by the correct GUT model. Like for quarks and charged leptons one expects a mass hierarchy also for the right-handed neutrinos. For instance, if their masses scale like the up-quark masses one has $M_1 \sim 10^{-5} M_3 \sim 10^{10}$ GeV.

The lightest of the heavy Majorana neutrinos, $N_1$, is ideally suited to generate the cosmological baryon asymmetry. Since it has no standard model gauge interactions it can naturally satisfy the out-of-equilibrium condition. $N_1$ decays to lepton-Higgs pairs then yield a lepton asymmetry $\langle L \rangle_T \neq 0$, which is partially converted to a baryon asymmetry $\langle B \rangle_T \neq 0$. The generated asymmetry is proportional to the $CP$ asymmetry in $N_1$-decays. In the case of the standard model with one Higgs doublet, i.e. $H_1 = H_2^* = \phi$, the $CP$ asymmetry is conveniently written in the following form,

$$
\varepsilon_1 = \frac{\Gamma(N_1 \rightarrow l\phi) - \Gamma(N_1 \rightarrow \bar{l}\bar{\phi})}{\Gamma(N_1 \rightarrow l\phi) + \Gamma(N_1 \rightarrow \bar{l}\bar{\phi})} \approx \frac{3}{16\pi} \frac{M_1}{(h^\dagger \nu h^\nu)_{11} v^2} \text{Im} (h^\dagger \nu m_\nu h^\nu)_{11}, \quad (14)
$$

where the seesaw mass relation (13) has been used.

From the expression (14) one easily obtains a rough estimate for $\varepsilon_1$ in terms of neutrino masses. Assuming dominance of the largest eigenvalue of $m_\nu$, phases $O(1)$ and approximate cancellation of Yukawa couplings in numerator and denominator one finds,

$$
\varepsilon_1 \sim \frac{3}{16\pi} \frac{M_1 m_3}{v^2} \sim 0.1 \frac{M_1}{M_3}, \quad (15)
$$

where we have again used the seesaw relation. Hence, the order of magnitude of the $CP$ asymmetry is approximately given by the mass hierarchy of the heavy Majorana neutrinos. For $M_1 / M_3 \sim 10^{-5}$ one has $\varepsilon_1 \sim 10^{-6}$.

Given the $CP$ asymmetry $\varepsilon_1$ one obtains for the baryon asymmetry,

$$
\eta_B = \frac{n_B - n_{\bar{B}}}{n_\gamma} = \frac{\kappa}{f} c_S \varepsilon_1 \sim 10^{-9}. \quad (16)
$$

Here $f \sim 10^2$ is the dilution factor which accounts for the increase of the number of photons in a comoving volume element between baryogenesis and today. The
baryogenesis temperature is

$$T_B \sim M_1 \sim 10^{10} \text{ GeV},$$

(17)

and the washout factor $\kappa$ depends on the neutrino masses in a way, which will be discussed in detail in the following chapter. Its determination requires the solution of the Boltzmann equations\[^{19,20}\]; in the estimate (16) we have assumed a typical value, $\kappa \sim 0.1$. The correct value of the baryon asymmetry, $\eta_B \sim 10^{-9}$ is then obtained as consequence of a large hierarchy of the heavy neutrino masses, which leads to a small $CP$ asymmetry, and the kinematical factors $f$ and $\kappa^{22}$. The baryogenesis temperature $T_B \sim 10^{10} \text{ GeV}$ corresponds to the time $t_B \sim 10^{-26} \text{ s}$, which characterizes the next relevant epoch before recombination, nucleosynthesis and electroweak transition.

An important question concerns the relation between leptogenesis and neutrino mass matrices which can account for low-energy neutrino data\[^{23,24,25,26,27}\]. At present we know two mass differences for the light neutrinos and we have some information about elements of the mixing matrix $U$ in the leptonic charged current. Since $U$ could be entirely due to mixings among the charged leptons, this does not constrain the light neutrino mass matrix in a model independent way. The light neutrino masses can be either quasi-degenerate or hierarchical, and they can easily be consistent with successfull leptogenesis. In grand unified theories, due to quark-lepton unification, the hierarchical quark masses together with the seesaw relation require also hierarchical heavy Majorana masses. Neglecting the mixing of charged leptons, the constraints become very stringent and successful leptogenesis is possible only for a very small part of parameter space\[^{28}\]. In this case the enhancement of the $CP$ asymmetry for partially degenerate heavy neutrinos\[^{29,30,31}\] plays an important role. On the other hand, in unified theories large mixings between charged leptons easily occur. For instance, in a six-dimensional SO(10) model the parameters used in the above estimates for the baryon asymmetry were recently obtained\[^{32}\].

3. Quantitative Analysis and Bounds on Neutrino Masses

Leptogenesis is a non-equilibrium process which takes place at temperatures $T \sim M_1$. For a decay width small compared to the Hubble parameter, $\Gamma_1(T) < H(T)$, heavy neutrinos are out of thermal equilibrium, otherwise they are in thermal equilibrium. A rough estimate of the borderline between the two regimes is given by $\Gamma_1 = H(M_1)^{11}$. This is equivalent to the condition that the effective neutrino mass,

$$\tilde{m}_1 = \frac{(m_D^\dagger m_D)_{11}}{M_1},$$

(18)
equals the ‘equilibrium neutrino mass’

\[ m_* = \frac{16\pi^{5/2}}{3\sqrt{5}} g_*^{1/2} \frac{v^2}{M_{pl}} \simeq 10^{-3} \text{ eV}. \]  

(19)

Here we have used \( M_{pl} = 1.2 \times 10^{19} \text{ GeV} \) and \( g_* = 434/4 \) as effective number of degrees of freedom. For \( \tilde{m}_1 > m_* \) (\( \tilde{m}_1 < m_* \)) the heavy neutrinos of type \( N_1 \) are in (out of) thermal equilibrium at \( T = M_1 \).

It is very remarkable that the equilibrium neutrino mass \( m_* \) is close to the neutrino masses suggested by neutrino oscillations, \( \sqrt{\Delta m_{\text{sol}}^2} \simeq 8 \times 10^{-3} \text{ eV} \) and \( \sqrt{\Delta m_{\text{atm}}^2} \simeq 5 \times 10^{-2} \text{ eV} \). This suggests that it may be possible to understand the cosmological baryon asymmetry via leptogenesis as a process close to thermal equilibrium. Ideally, \( \Delta L = 1 \) and \( \Delta L = 2 \) processes would be strong enough at temperatures above \( M_1 \) to keep the heavy neutrinos in thermal equilibrium and weak enough to allow the generation of an asymmetry at temperatures below \( M_1 \).

In general, the generated baryon asymmetry is the result of a competition between production processes and washout processes which tend to erase any generated asymmetry. Unless the heavy Majorana neutrinos are partially degenerate, \( M_{2,3} - M_1 \ll M_1 \), the dominant processes are decays and inverse decays of \( N_1 \) and the usual off-shell \( \Delta L = 1 \) and \( \Delta L = 2 \) scatterings. The Boltzmann equations for leptogenesis read,

\[
\frac{dN_{N_1}}{dz} = -(D + S) (N_{N_1} - N_{N_1}^{\text{eq}}), \quad (20)
\]

\[
\frac{dN_{B-L}}{dz} = -\varepsilon_1 D (N_{N_1} - N_{N_1}^{\text{eq}}) - W N_{B-L}. \quad (21)
\]

Here \( N_i \) are number densities, \( z = M_1/T \) and \( D/(Hz) \), \( S/(Hz) \) and \( W/(Hz) \) denote decay rate, \( \Delta L = 1 \) scattering rate and \( \Delta L = 2 \) washout rate, respectively; all rates are normalized to the Hubble parameter.

In order to understand the dependence of the solutions on the neutrino parameters, it is crucial to note that the rates \( D/(Hz) \), \( S/(Hz) \) and the resonance contribution \( (W - \Delta W)/(Hz) \), are all proportional to the effective neutrino mass \( \tilde{m}_1 \). One finds,

\[
D, S, W - \Delta W \propto \frac{M_{pl}\tilde{m}_1}{v^2}, \quad \Delta W \propto \frac{M_{pl}M_1}{v^4} \bar{m}^2. \quad (22)
\]

Here \( \tilde{m}_1 \) is the effective neutrino mass \( (18) \), and \( \bar{m} \) is a quadratic mean,

\[
\bar{m}^2 = \text{tr} \left( m_\nu^\dagger m_\nu \right) = m_1^2 + m_2^2 + m_3^2. \quad (23)
\]

Eq. (22) implies that, as long as \( \Delta W \) can be neglected, the generated lepton asymmetry is independent of \( M_1 \). For quasi-degenerate neutrinos, with increasing \( \bar{m} \), the...
washout rate $\Delta W$ becomes important and eventually prevents successful leptogenesis. This leads to the upper bound on the absolute neutrino mass scale discussed below.

The decay, scattering and washout rates are shown in fig. 3 as functions of $z$ for a typical set of neutrino parameters, $M_1 = 10^{10}$ GeV, $\tilde{m}_1 = 10^{-3}$ eV, $\bar{m} = 0.05$ eV. All rates are of order the Hubble parameter at $z \sim 1$ where baryogenesis takes place. The generation of the $B-L$ asymmetry for these parameters is shown in fig. 4 for $|\varepsilon_1| = 10^{-6}$ and for two different initial conditions: zero and thermal $N_1$ abundance. The figure demonstrates that the Yukawa interactions are strong enough to bring the heavy neutrinos into thermal equilibrium before leptogenesis takes place. The resulting asymmetry is in accord with observation, $\eta_B \sim 0.01 \times N_{B-L} \sim 10^{-9}$.

Given the heavy neutrino mass $M_1$, the $CP$ asymmetry $\varepsilon_1$ satisfies an upper bound $|\varepsilon| \leq 10^{-7}$. Using this bound one can determine the maximal baryon asymmetry $\eta_{B}^{\text{max}}$ as function of the masses $\tilde{m}_1, M_1,$ and $\bar{m}$:

$$
\eta_B \leq \eta_{B}^{\text{max}}(\tilde{m}_1, M_1, \bar{m}) \simeq 0.96 \times 10^{-2} \varepsilon_1^{\text{max}}(\tilde{m}_1, M_1, \bar{m}) \kappa(\tilde{m}_1, M_1, \bar{m}).
$$

(24)

Requiring the maximal baryon asymmetry to be larger than the observed one,

$$
\eta_{B}^{\text{max}}(\tilde{m}_1, M_1, \bar{m}) \geq \eta_{B}^{\text{CMB}},
$$

(25)
yields a constraint on the neutrino mass parameters $\tilde{m}_1, M_1$ and $\overline{m}$.

The maximal $CP$ asymmetry as function of $\tilde{m}_1, M_1$ and $\overline{m}$ is given by

$$\varepsilon_{\text{max}}^1 = \frac{3}{16\pi} \frac{M_1 m_3}{v^2} \left[ 1 - \frac{m_1}{m_3} \left( 1 + \frac{m_3 - m_1}{\overline{m}_1} \right)^{1/2} \right].$$

(26)

For neutrino masses with normal hierarchy one has

$$m_3^2 - m_2^2 = \Delta m^2_{\text{atm}}, \quad m_2^2 - m_1^2 = \Delta m^2_{\text{sol}},$$

(27)

and the dependence on $\overline{m}$ reads

$$m_3^2 = \frac{1}{3} \left( \overline{m}^2 + 2\Delta m^2_{\text{atm}} + \Delta m^2_{\text{sol}} \right),$$

(28)

$$m_2^2 = \frac{1}{3} \left( \overline{m}^2 - \Delta m^2_{\text{atm}} + \Delta m^2_{\text{sol}} \right),$$

(29)

$$m_1^2 = \frac{1}{3} \left( \overline{m}^2 - \Delta m^2_{\text{atm}} - 2\Delta m^2_{\text{sol}} \right).$$

(30)

Note that $\varepsilon_1 = 0$ for $\overline{m}_1 = m_1$. In general, one expects $m_1 \leq \overline{m}_1 \lesssim m_3$. Here the lower bound is always true whereas the upper bounds holds if there are no strong
cancellations due to phase relations between different elements of the neutrino mass matrix.

Using the upper bound (26) on the CP asymmetry one can calculate the maximal baryon asymmetry. The CMB constraint (25) yields for each value of $\mu$ a domain in the $(\tilde{m}_1-M_1)$-plane which is allowed by successful baryogenesis (cf. fig. 5). For $\mu < 0.20$ eV this domain shrinks to zero. Using the relations (28)-(30), one can easily translate this bound into upper limits on the individual neutrino masses,

$$m_1, m_2 < 0.11 \text{ eV} , \quad m_3 < 0.12 \text{ eV} . \quad (31)$$

Note that these bounds are a factor of factor of two below the recent upper bound of 0.23 eV obtained by WMAP.[2] For an inverted hierarchy of neutrino masses one finds very similar upper bounds, $m_1 < 0.11$ eV and $m_2, m_3 < 0.12$ eV. In a complete analysis the change of neutrino masses between the mass scale of leptogenesis and the electroweak scale has to be included. Generically, one expects that this will make the upper bounds on the neutrino masses more stringent.[38]

In a similar way one can obtain a lower bound on $M_1$, the smallest mass of the heavy Majorana neutrinos. One finds

$$M_1 > 4 \times 10^8 \text{ GeV} . \quad (32)$$
As a consequence, thermal leptogenesis requires a rather high reheating temperature, $T_R \gtrsim T_B \sim M_1$.

An important question for leptogenesis, and baryogenesis in general, is the dependence on initial conditions. This includes the dependence on the initial abundance of heavy Majorana neutrinos, as discussed above, and also the effect of an initial asymmetry which may have been generated by some other mechanism. It turns out that heavy Majorana neutrinos can efficiently erase initial asymmetries.

For large initial asymmetries, one can neglect the small asymmetry generated through the $CP$ violating interactions of the heavy neutrinos, i.e. one may set $\varepsilon_1 = 0$. The kinetic equation (21) for the asymmetry then becomes

$$\frac{dN_{B-L}}{dz} = -W N_{B-L} , \quad (33)$$

where $-N_{B-L}$ is the number of lepton doublets per comoving volume. The final $B-L$ asymmetry is then given by

$$N_{B-L}^f = \varphi(z_i) N_{B-L}^i , \quad (34)$$

with the washout factor

$$\varphi(z_i) = e^{-\int_{z_i}^{\infty} dz W(z)} . \quad (35)$$
The result of a quantitative analysis is shown in fig. 6. The washout becomes very efficient for $\tilde{m}_1 > m_* \simeq 10^{-3}$ eV. Already for $\tilde{m}_1 = 5 \times 10^{-3}$ eV one has $w(z_i) \sim 10^{-7} \ldots 10^{-5}$, indicated by the dashed and full lines, respectively. Hence, an initial asymmetry several orders of magnitude larger than the presently observed one can be reduced to a value below the one generated in leptogenesis. The range for $w(z_i)$ is due to a theoretical uncertainty in the treatment of $N_1$-top scatterings. Note, that a plateau for $w(z_i)$ is reached for values of $z_i$ just below one.

We conclude that leptogenesis naturally explains the observed baryon asymmetry for neutrino masses in the range

$$10^{-3} \text{ eV} \leq m_i \lesssim 0.1 \text{ eV} ,$$

almost independent of possible other pre-existing asymmetries. It is very remarkable that the data on solar and atmospheric neutrinos indicate neutrino masses precisely in this range.

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5. References

1) B. D. Fields, S. Sarkar, in Review of Particle Physics, Phys. Rev. D 66 (2002) 010001
2) WMAP Collaboration, D. N. Spergel et al., astro-ph/0302209
3) A. D. Sakharov, JETP Lett. 5 (1967) 24
4) E. W. Kolb, M. S. Turner, The Early Universe, Addison-Wesley, New York, 1990
5) For a recent review and references, see K. Hamaguchi, hep-ph/0212305
6) M. Fukugita, T. Yanagida, Phys. Lett. B 174 (1986) 45
7) G. ’tHooft, Phys. Rev. Lett. 37 (1976) 8; Phys. Rev. D 14 (1976) 3422
8) V. A. Kuzmin, V. A. Rubakov, M. E. Shaposhnikov, Phys. Lett. B 155 (1985) 36
9) D. Bödeker, G. D. Moore, K. Rummukainen, Phys. Rev. D 61 (2000) 056003
10) M. Quiros, Nucl. Phys. Proc. Suppl. 101 (2001) 401
11) K. Dick, M. Lindner, M. Ratz, D. Wright, Phys. Rev. Lett. 84 (2000) 4039
12) H. Murayama, A. Pierce, Phys. Rev. Lett. 89 (2002) 271601-1
13) H. Georgi, S. L. Glashow, Phys. Rev. Lett. 32 (1974) 438
14) H. Georgi, in Particles and Fields, ed. C. E. Carlson (AIP, NY, 1975) p. 575;
   H. Fritzsch, P. Minkowski, Ann. of Phys. 93 (1975) 193
15) T. Yanagida, in Workshop on Unified Theories, KEK report 79-18 (1979) p. 95;
   S. L. Glashow, in Quarks and Leptons, Cargèse 1979 (Plenum Press, New York, 1980) eds.
   M. Levy et al., p. 687;
   M. Gell-Mann, P. Ramond, R. Slansky, in Supergravity (North Holland, Amsterdam, 1979)
   eds. P. van Nieuwenhuizen, D. Freedman, p. 315
16) M. Flanz, E. A. Paschos, U. Sarkar, Phys. Lett. B 345 (1995) 248; Phys. Lett. B 384
    (1996) 487 (E)
17) L. Covi, E. Roulet, F. Vissani, Phys. Lett. B 384 (1996) 169
18) W. Buchmüller, M. Plümacher, Phys. Lett. B 431 (1998) 354
19) M. A. Luty, Phys. Rev. D 45 (1992) 455
20) M. Plümacher, Z. Phys. C 74 (1997) 549; Nucl. Phys. B 530 (1998) 207
21) R. Barbieri, P. Creminelli, A. Strumia, N. Tetradis, Nucl. Phys. B 575 (2000) 61;
    hep-ph/9911315-3
22) W. Buchmüller, M. Plümacher, Phys. Lett. B 389 (1996) 73
23) G. Altarelli, these proceedings
24) F. Buccella, these proceedings
25) F. Feruglio, these proceedings
26) S. King, these proceedings
27) R. N. Mohapatra, these proceedings
28) E. Kh. Akhmedov, M. Frigerio, A. Yu. Smirnov, hep-ph/0305322
29) A. Pilaftsis, Int. J. Mod. Phys. A 14 (1999) 1811
30) J. Ellis, M. Raidal, T. Yanagida, Phys. Lett. B 546 (2002) 228
31) G. C. Branco et al., Phys. Rev. D 67 (2003) 073025
32) T. Asaka, W. Buchmüller, L. Covi, hep-ph/0304142
33) W. Buchmüller, P. Di Bari, M. Plümacher, Nucl. Phys. B 643 (2002) 367
34) K. Hamaguchi, H. Murayama, T. Yanagida, Phys. Rev. D 65 (2002) 043512
35) S. Davidson, A. Ibarra, Phys. Lett. B 535 (2002) 25
36) W. Buchmüller, P. Di Bari, M. Plümacher, hep-ph/0302092
37) M. Fujii, K. Hamaguchi, T. Yanagida, Phys. Rev. D 65 (2002) 115012
38) S. Antusch, J. Kersten, M. Lindner, M. Ratz, hep-ph/0305273