MATTER ANTIMATTER ASYMMETRY
AND NEUTRINO PROPERTIES

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Abstract

The cosmological baryon asymmetry can be explained as remnant of heavy Majorana neutrino decays in the early universe. We study this mechanism for two models of neutrino masses with a large $\nu_\mu - \nu_\tau$ mixing angle which are based on the symmetries $SU(5) \times U(1)_F$ and $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_F$, respectively. In both cases $B - L$ is broken at the unification scale $\Lambda_{GUT}$. The models make different predictions for the baryogenesis temperature and the gravitino abundance.

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1 Baryogenesis and lepton number violation

The cosmological matter antimatter asymmetry, the ratio of the baryon density to the entropy density of the universe,

$$Y_B = \frac{n_B - n_{\bar{B}}}{s} = (0.6 - 1) \cdot 10^{-10},$$  \hspace{1cm} (1)

can in principle be understood in theories where baryon number, C and CP are not conserved [1]. The presently observed value of the baryon asymmetry is then explained as a consequence of the spectrum and interactions of elementary particles, together with the cosmological evolution.

A crucial ingredient of baryogenesis is the connection between baryon number ($B$) and lepton number ($L$) 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[2]. At zero temperature this has no observable effect due to the smallness of the weak coupling. However, as the temperature approaches the critical temperature $T_{EW}$ of the electroweak phase transition, $B$ and $L$ violating processes come into thermal equilibrium[3]. These ‘sphaleron processes’ violate baryon and lepton number by three units,

$$\Delta B = \Delta L = 3.$$  \hspace{1cm} (2)

It is generally believed that $B$ and $L$ changing 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} (3)

The non-conservation of baryon and lepton number has a profound effect on the generation of the cosmological baryon asymmetry. Eq. 2 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 $Y_B$ and the corresponding $L$ and $B - L$ asymmetries $Y_L$ and $Y_{B-L}$, respectively[4],

$$Y_B = a \ Y_{B-L} = \frac{a}{a - 1} \ Y_L,$$  \hspace{1cm} (4)
where $a$ is a number $\mathcal{O}(1)$. In the standard model with three generations and two Higgs doublets one has $a = 8/23$.

We conclude that $B - L$ violation is needed if the baryon asymmetry is generated before the electroweak transition, i.e. at temperatures $T > T_{EW} \sim 100$ 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.

The remnant of lepton number violation at low energies is an effective $\Delta L = 2$ interaction between lepton and Higgs fields,

\begin{equation}
\mathcal{L}_{\Delta L=2} = \frac{1}{2} f_{ij} l_i^T H_2 C l_j H_2 + \text{h.c.} .
\end{equation}

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 $\langle H_2 \rangle = v_2$, it gives rise to Majorana masses of the light neutrinos $\nu_e, \nu_\mu$ and $\nu_\tau$.

At finite temperature the $\Delta L = 2$ processes described by (5) take place with the rate\cite{5}

\begin{equation}
\Gamma_{\Delta L=2}(T) = \frac{1}{\pi^3} \frac{T^3}{v_2^4} \sum_{i=e,\mu,\tau} m_{\nu_i}^2 .
\end{equation}

In thermal equilibrium this yields an additional relation between the chemical potentials which implies

\begin{equation}
Y_B = Y_{B-L} = Y_L = 0 .
\end{equation}

To avoid this conclusion, the $\Delta L = 2$ interaction (5) must not reach thermal equilibrium. For baryogenesis at a temperature $T_B < T_{SPH} \sim 10^{12}$ GeV, one has to require $\Gamma_{\Delta L=2} < H|_{T_B}$, where $H$ is the Hubble parameter. This yields a stringent upper bound on Majorana neutrino masses,

\begin{equation}
\sum_{i=e,\mu,\tau} m_{\nu_i}^2 < \left( \frac{0.2 \text{ eV}}{T_{SPH}/T_B} \right)^{1/2} .
\end{equation}

For $T_B \sim T_{SPH}$, this bound would be comparable to the upper bound on the electron neutrino mass obtained from neutrinoless double beta decay. However, eq. (8) also applies to the $\tau$-neutrino mass. Note, that the bound can
be evaded if appropriate asymmetries are present for particles which reach thermal equilibrium only at temperatures below $T_B$ [6].

The connection between lepton number and the baryon asymmetry is lost if baryogenesis takes place at or below the Fermi scale[7]. However, detailed studies of the thermodynamics of the electroweak transition have shown that, at least in the standard model, the deviation from thermal equilibrium is not sufficient for baryogenesis[8]. In the minimal supersymmetric extension of the standard model (MSSM) such a scenario appears still possible for a limited range of parameters[7].

## 2 Decays of heavy Majorana neutrinos

Baryogenesis above the Fermi scale requires $B - L$ violation, and therefore $L$ violation. Lepton number violation is most simply realized by adding right-handed Majorana neutrinos to the standard model. Heavy right-handed Majorana neutrinos, whose existence is predicted by all extensions of the standard model containing $B - L$ as a local symmetry, can also explain the smallness of the light neutrino masses via the see-saw mechanism[9].

The most general Lagrangian for couplings and masses of charged leptons and neutrinos reads

$$\mathcal{L}_Y = -h_{eij}\bar{e}_R^i l_j H_1 - h_{\nu ij}\bar{\nu}_R^i l_j H_2 - \frac{1}{2}h_{\nu ij}\bar{\nu}_R^i \nu_R^j R + \text{h.c.} .$$

(9)

The vacuum expectation values of the Higgs field $\langle H_1 \rangle = v_1$ and $\langle H_2 \rangle = v_2 = \tan \beta \ v_1$ generate Dirac masses $m_l$ and $m_D$ for charged leptons and neutrinos, $m_e = h_e v_1$ and $m_D = h_\nu v_2$, respectively, which are assumed to be much smaller than the Majorana masses $M = h_r \langle R \rangle$. This yields light and heavy neutrino mass eigenstates

$$\nu \simeq K^\dagger \nu_L + \nu_R^c K , \quad N \simeq \nu_R + \nu_R^c ,$$

(10)

with masses

$$m_\nu \simeq -K^\dagger m_D \frac{1}{M} m_D^T K^* , \quad m_N \simeq M .$$

(11)

Here $K$ is a unitary matrix which relates weak and mass eigenstates.

The right-handed neutrinos, whose exchange may erase any lepton asymmetry, can also generate a lepton asymmetry by means of out-of-equilibrium decays.
This lepton asymmetry is then partially transformed into a baryon asymmetry by sphaleron processes\cite{10}. The decay width of the heavy neutrino $N_i$ reads at tree level,

$$
\Gamma_{Di} = \Gamma (N_i \rightarrow H^2_{2} + l) + \Gamma (N_i \rightarrow H^2_{2} + l^c) = \frac{1}{8\pi} (h^\nu_i h^\nu_i)_{ii} M_i .
$$

(12)

From the decay width one obtains an upper bound on the light neutrino masses via the out-of-equilibrium condition\cite{11}. Requiring $\Gamma_{D1} < H|_{T=M_1}$ yields the constraint

$$
\bar{m}_1 = (h^\nu_i h^\nu_i)_{11} \frac{v^2_2}{M_1} < 10^{-3} \text{eV} .
$$

(13)

More direct bounds on the light neutrino masses depend on the structure of the Dirac neutrino mass matrix.

Interference between the tree-level amplitude and the one-loop self-energy and vertex corrections yields $CP$ asymmetries in the heavy Majorana neutrino decays. In a basis, where the right-handed neutrino mass matrix $M = h_r \langle R \rangle$ is diagonal, one obtains\cite{12,13}

$$
\varepsilon_1 = \frac{\Gamma(N_1 \rightarrow l H^2_2) - \Gamma(N_1 \rightarrow l^c H^2_{2})}{\Gamma(N_1 \rightarrow l H^2_2) + \Gamma(N_1 \rightarrow l^c H^2_{2})} \\
\simeq -\frac{3}{16\pi} \frac{1}{(h^\nu \nu^\dagger)_{11}} \sum_{i=2,3} \text{Im} \left[ (h^\nu_i h^\nu_i)_{11}^2 \frac{M_i}{M_1} \right] .
$$

(14)

Here we have assumed $M_1 < M_2, M_3$, which is satisfied in the applications considered in the following sections.

In the early universe at temperatures $T \sim M_1$ the CP asymmetry (14) leads to a lepton asymmetry\cite{14},

$$
Y_L = \frac{n_L - n_T}{s} = \kappa \frac{\varepsilon_1}{g_*} .
$$

(15)

Here the factor $\kappa < 1$ represents the effect of washout processes. In order to determine $\kappa$ one has to solve the full Boltzmann equations \cite{15,16}. In the examples discussed below one has $\kappa \simeq 10^{-1} \ldots 10^{-3}$. 


3 Neutrino masses and mixings

The CP asymmetry (14) is given in terms of the Dirac and the Majorana neutrino mass matrices. Depending on the neutrino mass hierarchy and the size of the mixing angles the CP asymmetry can vary over many orders of magnitude. It is therefore interesting to see whether a pattern of neutrino masses motivated by other considerations is consistent with leptogenesis.

An attractive framework to explain the observed mass hierarchies of quarks and charged leptons is the Froggatt-Nielsen mechanism [18] based on a spontaneously broken $U(1)_F$ generation symmetry. The Yukawa couplings arise from non-renormalizable interactions after a gauge singlet field $\Phi$ acquires a vacuum expectation value,

$$h_{ij} = g_{ij} \left( \frac{\langle \Phi \rangle}{\Lambda} \right)^{Q_i + Q_j} \ .$$

Here $g_{ij}$ are couplings $O(1)$ and $Q_i$ are the $U(1)$ charges of the various fermions with $Q_\Phi = -1$. The interaction scale $\Lambda$ is expected to be very large, $\Lambda > \Lambda_{GUT}$. In the following we shall discuss two different realizations of this idea which are motivated by the recently reported atmospheric neutrino anomaly [17]. Both scenarios have a large $\nu_\mu - \nu_\tau$ mixing angle. They differ, however, by the symmetry structure and by the size of the parameter $\epsilon$ which characterizes the flavour mixing.

3.1 $SU(5) \times U(1)_F$

This symmetry has been considered by a number of authors [22]. Particularly interesting is the case with a nonparallel family structure where the chiral $U(1)_F$ charges are different for the $5^*$-plets and the $10$-plets of the same family [19]-[21]. An example of possible charges $Q_i$ is given in table 1.

The assignment of the same charge to the lepton doublets of the second and

| $\psi_i$ | $e_{R3}^c$ | $e_{R2}^c$ | $e_{R1}^c$ | $l_{L3}$ | $l_{L2}$ | $l_{L1}$ | $\nu_{R3}^c$ | $\nu_{R2}^c$ | $\nu_{R1}^c$ |
|----------|-----------|-----------|-----------|--------|--------|--------|-----------|-----------|-----------|
| $Q_1$    | 0         | 1         | 2         | 0      | 0      | 1      | 0         | 1         | 2         |

Table 1

Chiral charges of charged and neutral leptons with $SU(5) \times U(1)_F$ symmetry [23].
third generation leads to a neutrino mass matrix of the form [19,20],

$$m_{\nu ij} \sim \begin{pmatrix} e^2 & \epsilon & \epsilon \\ \epsilon & 1 & 1 \\ \epsilon & 1 & 1 \end{pmatrix} \frac{v_2^2}{\langle R \rangle}.$$  \hspace{1cm} (17)

This structure immediately yields a large $\nu_\mu - \nu_\tau$ mixing angle. The phenomenology of neutrino oscillations depends on the unspecified coefficients $O(1)$. The parameter $\epsilon$ which gives the flavour mixing is chosen to be

$$\langle \Phi \rangle_\Lambda = \epsilon \sim \frac{1}{17}.$$  \hspace{1cm} (18)

The three Yukawa matrices for the leptons have the structure,

$$h_e \sim \begin{pmatrix} e^3 & e^2 & e^2 \\ e^2 & \epsilon & \epsilon \\ \epsilon & 1 & 1 \end{pmatrix}, \quad h_\nu \sim \begin{pmatrix} e^3 & e^2 & e^2 \\ e^2 & \epsilon & \epsilon \\ \epsilon & 1 & 1 \end{pmatrix}, \quad h_r \sim \begin{pmatrix} e^4 & e^3 & e^2 \\ e^3 & e^2 & \epsilon \\ e^2 & \epsilon & 1 \end{pmatrix}.$$  \hspace{1cm} (19)

Note, that $h_e$ and $h_\nu$ have the same, non-symmetric structure. One easily verifies that the mass ratios for charged leptons, heavy and light Majorana neutrinos are given by

$$m_e : m_\mu : m_\tau \sim \epsilon^3 : \epsilon : 1, \quad M_1 : M_2 : M_3 \sim \epsilon^4 : \epsilon^2 : 1,$$

$$m_1 : m_2 : m_3 \sim \epsilon^2 : 1 : 1.$$  \hspace{1cm} (20) \hspace{1cm} (21)

The masses of the two eigenstates $\nu_\mu$ and $\nu_\tau$ depend on unspecified factors of order one, and may easily differ by an order of magnitude [25]. They can therefore be consistent with the mass differences $\Delta m_{\nu_\mu\nu_\mu}^2 \simeq 4 \cdot 10^{-6} - 1 \cdot 10^{-5}$ eV$^2$ [27] inferred from the MSW solution of the solar neutrino problem [28] and $\Delta m_{\nu_\mu\nu_\tau}^2 \simeq (5 \cdot 10^{-4} - 6 \cdot 10^{-3})$ eV$^2$ associated with the atmospheric neutrino deficit [17]. For numerical estimates we shall use the average of the neutrino masses of the second and third family, $m_\nu = (m_{\nu_\mu} m_{\nu_\tau})^{1/2} \sim 10^{-2}$ eV. Note, that for a different choice of U(1) charges the coefficients in eq. (17) automatically yield the hierarchy $m_2/m_3 \sim \epsilon^{2/3}$ [29].

| $\psi_i$ | $e^c_{R3}$ | $e^c_{R2}$ | $e^c_{R1}$ | $l^c_{L3}$ | $l^c_{L2}$ | $l^c_{L1}$ | $\nu^c_{R3}$ | $\nu^c_{R2}$ | $\nu^c_{R1}$ |
|---------|-----------|-----------|-----------|-----------|-----------|-----------|-----------|-----------|-----------|
| $Q_i$   | $0$       | $\frac{1}{2}$ | $\frac{5}{2}$ | $0$       | $\frac{1}{2}$ | $\frac{5}{2}$ | $0$       | $\frac{1}{2}$ | $\frac{5}{2}$ |

Table 2

*Chiral charges of charged and neutral leptons with SU(3)$_c \times$ SU(3)$_L \times$ SU(3)$_R \times$ U(1)$_F$ symmetry [22].*
The choice of the charges in table 1 corresponds to large Yukawa couplings of the third generation. For the mass of the heaviest Majorana neutrino one finds

\[ M_3 \sim \frac{v^2}{m_\nu} \sim 10^{15} \text{ GeV}. \]  

(22)

This implies that \( B - L \) is broken at the unification scale \( \Lambda_{GUT} \).

### 3.2 \( SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_F \)

This symmetry arises in unified theories based on the gauge group \( E_6 \). The leptons \( e_R, l_L \) and \( \nu_R \) are contained in a single \((1,3,\bar{3})\) representation. Hence, all leptons of the same generation have the same \( U(1)_F \) charge and all leptonic Yukawa matrices are symmetric. Masses and mixings of quarks and charged leptons can be successfully described by using the charges given in table 2 [22]. Clearly, the three Yukawa matrices have the same structure\(^1\).

\[
\begin{align*}
\begin{pmatrix}
\epsilon^5 & \epsilon^3 & \epsilon^{5/2} \\
\epsilon^3 & \epsilon & \epsilon^{1/2} \\
\epsilon^{5/2} & \epsilon^{1/2} & 1
\end{pmatrix}, & \quad h_e, h_r \sim \begin{pmatrix}
\bar{\epsilon}^5 & \bar{\epsilon}^3 & \bar{\epsilon}^{5/2} \\
\bar{\epsilon}^3 & \bar{\epsilon} & \bar{\epsilon}^{1/2} \\
\bar{\epsilon}^{5/2} & \bar{\epsilon}^{1/2} & 1
\end{pmatrix}. \\

h_\nu \sim & \begin{pmatrix}
\epsilon^5 & \epsilon^3 & \epsilon^{5/2} \\
\epsilon^3 & \epsilon & \epsilon^{1/2} \\
\epsilon^{5/2} & \epsilon^{1/2} & 1
\end{pmatrix}.
\end{align*}
\]  

(23)

Note, that the expansion parameter in \( h_\nu \) is different from the one in \( h_e \) and \( h_r \). From the quark masses, which also contain \( \epsilon \) and \( \bar{\epsilon} \), one infers \( \bar{\epsilon} \simeq \epsilon^2 \) [22].

From eq. (23) one obtains for the masses of charged leptons, light and heavy Majorana neutrinos,

\[
\begin{align*}
\begin{array}{c}
m_e : m_\mu : m_\tau \\ m_1 : m_2 : m_3
\end{array} & \sim \begin{array}{c}
M_1 : M_2 : M_3 \\
\epsilon^5 : \epsilon : 1
\end{array} , \\
\begin{array}{c}
m_1 : m_2 : m_3 \\
\epsilon^{15} : \epsilon^3 : 1
\end{array}.
\end{align*}
\]  

(24)  

(25)

Like in the example with \( SU(5) \times U(1)_F \) symmetry, the mass of the heaviest Majorana neutrino,

\[
M_3 \sim \frac{v^2}{m_\nu} \sim 10^{15} \text{ GeV},
\]  

(26)

implies that \( B - L \) is broken at the unification scale \( \Lambda_{GUT} \).

\(^1\) Note, that with respect to ref. [22], \( \epsilon \) and \( \bar{\epsilon} \) have been interchanged.
The $\nu_\mu - \nu_\tau$ mixing angle is mostly given by the mixing of the charged leptons of the second and third generation [22],

$$\sin \Theta_{\mu\tau} \sim \sqrt{\epsilon + \epsilon}.$$  \hspace{1cm} (27)

This requires large flavour mixing,

$$\left( \frac{\langle \Phi \rangle}{\Lambda} \right)^{1/2} = \sqrt{\epsilon} \sim \frac{1}{2}.$$  \hspace{1cm} (28)

In view of the unknown coefficients $O(1)$ the corresponding mixing angle $\sin \Theta_{\mu\tau} \sim 0.7$ is consistent with the interpretation of the atmospheric neutrino anomaly as $\nu_\mu - \nu_\tau$ oscillation.

It is very instructive to compare the two scenarios of lepton masses and mixings described above. In the first case, the large $\nu_\mu - \nu_\tau$ mixing angle follows from a nonparallel flavour symmetry. The parameter $\epsilon$, which characterizes the flavour mixing, is small. In the second case, the large $\nu_\mu - \nu_\tau$ mixing angle is a consequence of the large flavour mixing $\epsilon$. The $U(1)_F$ charges of all leptons are the same, i.e., one has a parallel family structure. Also the mass hierarchies, given in terms of $\epsilon$, are rather different. This illustrates that the separation into a flavour mixing parameter $\epsilon$ and coefficients $O(1)$ is far from unique. It is therefore important to study other observables which depend on the lepton mass matrices. A particular example is the baryon asymmetry.

4 Matter antimatter asymmetry

We can now evaluate the baryon asymmetry for the two patterns of neutrino mass matrices discussed in the previous section. A rough estimate of the baryon asymmetry can be obtained from the CP asymmetry $\varepsilon_1$ of the heavy Majorana neutrino $N_1$. A quantitative determination requires a numerical study of the full Boltzmann equations [16].

4.1 $SU(5) \times U(1)_F$

In this case one obtains from eqs. (14) and (19),

$$\varepsilon_1 \sim \frac{3}{16\pi} \epsilon^4.$$  \hspace{1cm} (29)
Fig. 1. Time evolution of the neutrino number density and the lepton asymmetry in the case of the $SU(5) \times U(1)_F$ symmetry. The solid line shows the solution of the Boltzmann equation for the right-handed neutrinos, while the corresponding equilibrium distribution is represented by the dashed line. The absolute value of the lepton asymmetry $Y_L$ is given by the dotted line and the hatched area shows the lepton asymmetry corresponding to the observed baryon asymmetry.

From eq. (15), $\epsilon^2 \sim 1/300$ (18) and $g_* \sim 100$ one then obtains the baryon asymmetry,

$$Y_B \sim \kappa \ 10^{-8}.$$  \hfill (30)

For $\kappa \sim 0.1 \ldots 0.01$ this is indeed the correct order of magnitude. The baryogenesis temperature is given by the mass of the lightest of the heavy Majorana neutrinos,

$$T_B \sim M_1 \sim \epsilon^4 M_3 \sim 10^{10} \text{ GeV}.$$  \hfill (31)

This set of parameters, where the CP asymmetry is given in terms of the mass hierarchy of the heavy neutrinos, has been studied in detail [26]. The generated baryon asymmetry does not depend on the flavour mixing of the light neutrinos. The $\nu_\mu - \nu_\tau$ mixing angle is large in the scenario described in the previous section whereas it was assumed to be small in [26].

The solution of the full Boltzmann equations is shown in fig. 1 for the non-supersymmetric case [26]. The initial condition at a temperature $T \sim 10 M_1$ is chosen to be a state without heavy neutrinos. The Yukawa interactions are sufficient to bring the heavy neutrinos into thermal equilibrium. At temper-
Fig. 2. Solution of the Boltzmann equations in the case of the $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_F$ symmetry.

At temperatures $T \sim M_1$ this is followed by the usual out-of-equilibrium decays which lead to a non-vanishing baryon asymmetry. The final asymmetry agrees with the estimate (30) for $\kappa \sim 0.1$.

The change of sign in the lepton asymmetry is due to the fact that inverse decay processes, which take part in producing the neutrinos, are CP violating, i.e. they generate a lepton asymmetry at high temperatures. Due to the interplay of inverse decay processes and lepton number violating scattering processes this asymmetry has a different sign than the one produced by neutrino decays at lower temperatures.

4.2 $SU(3)_c \times SU(3)_L \times SU(3)_R \times U(1)_F$

In this case the neutrino Yukawa couplings (23) yield the CP asymmetry

$$\varepsilon_1 \sim \frac{3}{16\pi} \epsilon^5,$$  \hspace{1cm} (32)

which correspond to the baryon asymmetry (cf. (15))

$$Y_B \sim \kappa \ 10^{-6}.$$  \hspace{1cm} (33)

Due to the large value of $\epsilon$ the CP asymmetry is two orders of magnitude larger than in the case with $SU(5) \times U(1)_F$ symmetry. However, washout processes
are now also stronger. The solution of the Boltzmann equations is shown in fig. 2. The final asymmetry is again $Y_B \sim 10^{-9}$ which now corresponds to $\kappa \sim 10^{-3}$. The baryogenesis temperature is considerably larger than in the first case,

$$T_B \sim M_1 \sim e^5 M_3 \sim 10^{12} \text{ GeV}.$$  \hfill (34)

The baryon asymmetry is largely determined by the parameter $\tilde{m}_1$ defined in eq. (13) [16]. In the first example, one has $\tilde{m}_1 \sim \overline{m}_\nu$. In the second case one finds $\tilde{m}_1 \sim m_3$. Since $\overline{m}_\nu$ and $m_3$ are rather similar it is not too surprising that the generated baryon asymmetry is about the same in both cases.

5 Conclusions

Detailed studies of the thermodynamics of the electroweak interactions at high temperatures have shown that in the standard model and most of its extensions the electroweak transition is too weak to affect the cosmological baryon asymmetry. Hence, one has to search for baryogenesis mechanisms above the Fermi scale.

Due to sphaleron processes baryon number and lepton number are related in the high-temperature, symmetric phase of the standard model. As a consequence, the cosmological baryon asymmetry is related to neutrino properties. Baryogenesis requires lepton number violation, which occurs in extensions of the standard model with right-handed neutrinos and Majorana neutrino masses.

Although lepton number violation is needed in order to obtain a baryon asymmetry, it must not be too strong since otherwise any baryon and lepton asymmetry would be washed out. This leads to stringent upper bounds on neutrino masses which depend on the particle content of the theory.

The solar and atmospheric neutrino deficits can be interpreted as a result of neutrino oscillations. For hierarchical neutrinos the corresponding neutrino masses are very small. Assuming the see-saw mechanism, this suggests the existence of very heavy right-handed neutrinos and a large scale of $B - L$ breaking.

It is remarkable that these hints on the nature of lepton number violation fit very well together with the idea of leptogenesis. For hierarchical neutrino masses, with $B - L$ broken at the unification scale $\Lambda_{\text{GUT}} \sim 10^{16} \text{ GeV}$, the observed baryon asymmetry $Y_B \sim 10^{-10}$ is naturally explained by the decay of heavy Majorana neutrinos.
Although the observed baryon asymmetry imposes important constraints on neutrino properties, other observables are needed to discriminate between different models. The two examples considered in this paper predict different baryogenesis temperatures. Correspondingly, in supersymmetric models the predictions for the gravitino abundance are different [30]-[32]. In the case with $SU(5) \times U(1)_F$ symmetry, stable gravitinos can be the dominant component of cold dark matter [32]. The models make also different predictions for the rate of lepton flavour changing radiative corrections.

References

[1] A. D. Sakharov, JETP Lett. 5 (1967) 24
[2] G. ’t Hooft, Phys. Rev. Lett. 37 (1976) 8
[3] V. A. Kuzmin, V. A. Rubakov, M. E. Shaposhnikov, Phys. Lett. B 155 (1985) 36
[4] J. A. Harvey, M. S. Turner, Phys. Rev. D 42 (1990) 3344
[5] M. Fukugita, T. Yanagida, Phys. Rev. D 42 (1990) 1285
[6] J. M. Cline, K. Kainulainen, K. A. Olive, Phys. Rev. Lett. 71 (1993) 2372
[7] For a review and references, see
  A. D. Dolgov, Phys. Rep. 222C (1992) 309;
  V. A. Rubakov, M. E. Shaposhnikov, Phys. Usp. 39 (1996) 461;
  S. J. Huber, M. G. Schmidt, SUSY Variants of the Electroweak Phase Transition, hep-ph/9809506
[8] For a discussion and references, see
  K. Jansen, Nucl. Phys. B (Proc. Supp.) 47 (1996) 196;
  W. Buchmüller, in Quarks ’96 (Yaroslavl, Russia, 1996) eds. V. A. Matveev et al., hep-ph/9610335;
  K. Rummukainen, Nucl. Phys. B (Proc. Suppl.) 53 (1997) 30
[9] T. Yanagida, in Workshop on unified Theories, KEK report 79-18 (1979) p. 95;
  M. Gell-Mann, P. Ramond, R. Slansky, in Supergravity (North Holland, Amsterdam, 1979) eds. P. van Nieuwenhuizen, D. Freedman, p. 315
[10] M. Fukugita, T. Yanagida, Phys. Lett. B 174 (1986) 45
[11] W. Fischler, G. F. Giudice, R. G. Leigh, S. Paban, Phys. Lett. B 258 (1991) 45
[12] L. Covi, E. Roulet, F. Vissani, Phys. Lett. B 384 (1996) 169;
    M. Flanz, E. A. Paschos, U. Sarkar, Phys. Lett. B 345 (1995) 248; Phys. Lett. B 384 (1996) 487 (E)
[13] W. Buchmüller, M. Plüümacher, Phys. Lett. B 431 (1998) 354

[14] A. D. Dolgov, Ya. B. Zeldovich, Rev. Mod. Phys. 53 (1981) 1; E. W. Kolb, S. Wolfram, Nucl. Phys. B 172 (1980) 224; Nucl. Phys. B 195 (1982) 542(E)

[15] M. A. Luty, Phys. Rev. D 45 (1992) 455

[16] M. Plüümacher, Z. Phys. C 74 (1997) 549; Nucl. Phys. B 530 (1998) 207

[17] Super-Kamiokande Collaboration, Y. Fukuda et al., Phys. Rev. Lett. 81 (1998) 1562

[18] C. D. Froggatt, H. B. Nielsen, Nucl. Phys. B 147 (1979) 277

[19] J. Sato, T. Yanagida, Talk at Neutrino'98, hep-ph/9809307

[20] P. Ramond, Talk at Neutrino'98, hep-ph/9809401

[21] J. Bijnens, C. Wetterich, Nucl. Phys. B 292 (1987) 443

[22] For a recent discussion and references, see S. Lola, G. G. Ross, hep-ph/9902283

[23] W. Buchmüller, T. Yanagida, Phys. Lett. B 445 (1999) 399

[24] F. Vissani, JHEP11 (1998) 025

[25] N. Irges, S. Lavignac, P. Ramond, Phys. Rev. D 58 (1998) 035003

[26] W. Buchmüller, M. Plüümacher, Phys. Lett. B 389 (1996) 73

[27] N. Hata, P. Langacker, Phys. Rev. D 56 (1997) 6107

[28] S. P. Mikheyev, A. Y. Smirnov, Nuovo Cim. 9C (1986) 17; L. Wolfenstein, Phys. Rev. D 17 (1978) 2369

[29] G. Altarelli, F. Feruglio, JHEP11 (1998) 021; hep-ph/9812475

[30] M. Yu. Khlopov, A. D. Linde, Phys. Lett. B 138 (1984) 265; J. Ellis, J. E. Kim, D. V. Nanopoulos, Phys. Lett. B 145 (1984) 181

[31] T. Moroi, H. Murayama, M. Yamaguchi, Phys. Lett. B 303 (1993) 289

[32] M. Bolz, W. Buchmüller, M. Plüümacher, Phys. Lett. B 443 (1998) 209