Baryon Asymmetry, Lepton Mixing and SO(10) Unification *

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Abstract

Baryogenesis appears to require lepton number violation. This is naturally realized in extensions of the standard model containing right-handed neutrinos. We discuss the generation of a baryon asymmetry by the out-of-equilibrium decay of heavy Majorana neutrinos in these models, without and with supersymmetry. All relevant lepton number violating scattering processes which can inhibit the generation of an asymmetry are taken into account. We assume a similar pattern of mixings and masses for neutrinos and up-type quarks, as suggested by SO(10) unification. This implies that $B - L$ is broken at the unification scale $\Lambda_{\text{GUT}} \sim 10^{16}$ GeV, if $m_{\nu_\mu} \sim 3 \cdot 10^{-3}$ eV, as preferred by the MSW solution to the solar neutrino deficit. The observed baryon asymmetry is then obtained without any fine tuning of parameters.

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Figure 1: Contributions to the decay of a heavy Majorana neutrino

1. Standard model with right-handed neutrinos

Baryon number \((B)\) and lepton number \((L)\) are not conserved in the standard model. At temperatures above the critical temperature of the electroweak phase transition \((B + L)\) violating sphaleron processes are in thermal equilibrium. Hence, the cosmological baryon asymmetry appears to require \(B - L\) violation, and therefore \(L\) violation. Lepton number violation is naturally realized by adding right-handed Majorana neutrinos to the standard model. Heavy right-handed Majorana neutrinos, whose existence is predicted by theories based on gauge groups containing SO(10), can also explain the smallness of the light neutrino masses via the see-saw mechanism.

The most general Lagrangian for couplings and masses of charged leptons and neutrinos is given by

\[
\mathcal{L}_Y = \bar{l} l H \lambda_i^* e_R + \bar{l} l \epsilon H \lambda_\nu^* \nu_R - \frac{1}{2} \bar{\nu}_R^c M \nu_R + \text{h.c.} ,
\]

The vacuum expectation value of the Higgs field \(\langle H^0 \rangle = v \neq 0\) generates Dirac masses \(m_l = \lambda_l v\) and \(m_\nu = \lambda_\nu v\) for charged leptons and neutrinos, which are assumed to be much smaller than the Majorana masses \(M\).

Out-of-equilibrium decays of right-handed neutrinos can generate a lepton asymmetry, which is then partially transformed into a baryon asymmetry by sphaleron processes. The decay width of \(N_i\) in its rest frame reads at tree level,

\[
\Gamma_{D_i} = \Gamma_{rs} \left( N^i \to l H \right) + \Gamma_{rs} \left( N^i \to \bar{l} H^\dagger \right) = \frac{(m^i_D m_D)_{ii}}{8\pi v^2} M_i .
\]

Interference between the tree-level amplitude and one-loop corrections (cf. fig. III) gives rise to a \(CP\) asymmetry,

\[
\varepsilon_i = \frac{\Gamma(N_i \to l H) - \Gamma(N_i \to \bar{l} H^\dagger)}{\Gamma(N_i \to l H) + \Gamma(N_i \to \bar{l} H^\dagger)} = \frac{1}{8\pi v^2} \sum_{j \neq i} \text{Im} \left( \frac{(m^j_D m_D)_{ij}}{(m^i_D m_D)_{ii}} \right) f \left( \frac{M_j^2}{M_i^2} \right) ,
\]

where

\[
f(x) = \sqrt{x} \left[ \frac{2 - x}{1 - x} - (1 + x) \ln \left( \frac{1 + x}{x} \right) \right] .
\]
right-handed neutrinos have to be numerous before decaying, i.e., they have to be in thermal equilibrium at high temperatures. The Yukawa interactions (1) are too weak to achieve this and additional interactions are therefore needed. Since right-handed neutrinos are a necessary ingredient of SO(10) unified theories, it is natural to consider leptogenesis within an extended gauge model, contained in a SO(10) GUT. The minimal extension of the standard model is based on the gauge group

\[ G = SU(3)_C \times SU(2)_L \times U(1)_{Y} \times U(1)_{Y'} \subset SO(10) \]  
(5)

Here \( U(1)_{Y'} \), and therefore \( B - L \), is spontaneously broken, and the breaking scale is related to the heavy neutrino masses. The additional neutral gauge boson \( Z' \) accounts for pair creation and annihilation processes and for flavour transitions between heavy neutrinos of different generations. For appropriately chosen parameters these processes generate an equilibrium distribution of heavy neutrinos at high temperatures.

Of crucial importance are the \( \Delta L = 2 \) lepton number violating scatterings shown in fig. 2 which, if too strong, erase any lepton asymmetry. Similarly, the \( \Delta L = 1 \) lepton number violating neutrino top-quark scatterings shown in fig. 3 have to be taken into account because of the large top Yukawa coupling. Finally, the heavy neutrino decays (cf. fig. 1) as well as the inverse decays have to be incorporated in the Boltzmann equations.

Based on these equations the resulting lepton and baryon asymmetries can be evaluated\(^7\),\(^6\), and one may ask whether the right order of magnitude of the asymmetry results naturally in the leptogenesis scenario. To address this question one has to discuss patterns of neutrino mass matrices which determine the generated asymmetry.

2. Neutrino masses and mixings

In the following we shall assume a similar pattern of mixings and mass ratios for leptons and quarks\(^8\), which is natural in SO(10) unification. Such an ansatz is most transparent in a basis where all mass matrices are maximally diagonal. In addition to real mass eigenvalues two mixing matrices then appear. One can always choose a basis for the lepton fields such that the mass matrices \( m_l \) for the charged leptons and \( M \) for the heavy Majorana neutrinos \( N_i \) are diagonal with real and positive eigenvalues. In this basis \( m_D \) is a general complex matrix, which can be diagonalized by a biunitary transformation. Therefore, we can write \( m_D \) as product of a diagonal matrix \( m_{D, \text{diag}} \) and two unitary matrices \( V \) and \( U^\dagger \),

\[ m_D = V m_{D, \text{diag}} U^\dagger \]  
(6)

![Figure 2: Lepton number violating lepton Higgs scattering](image-url)
where the eigenvalues $m_i$ of $m_{D,\text{diag}}$ are real and positive. In the absence of a Majorana mass term $V$ and $U$ would correspond to Kobayashi-Maskawa type mixing matrices of left- and right-handed charged currents, respectively.

According to eq. (3) the $CP$ asymmetry is determined by the mixings and phases present in the product $m_D^\dagger m_D$, where the matrix $V$ drops out. Hence, to leading order, the mixings and phases which are responsible for baryogenesis are entirely determined by the matrix $U$. Correspondingly, the mixing matrix in the leptonic charged current, which determines $CP$ violation and mixings of the light leptons, depends on mass ratios and mixing angles and phases of $U$ and $V$. This implies that there exists no direct connection between the $CP$ violation and generation mixing relevant at high and low energies.

Consider now the mixing matrix $U$. One can factor out five phases,

$$U = e^{i\gamma} e^{i\lambda_3 \alpha} e^{i\lambda_8 \beta} U_1 e^{i\lambda_3 \sigma} e^{i\lambda_8 \tau},$$

where the $\lambda_i$ are the Gell-Mann matrices. The remaining matrix $U_1$ depends on three mixing angles and one phase, like the CKM matrix for quarks. In analogy to the quark mixing matrix we choose the Wolfenstein parametrization\textsuperscript{[8]} as ansatz for $U_1$,

$$U_1 = \begin{pmatrix}
1-\lambda^2/2 & \lambda & A\lambda^3 (\rho - i\eta) \\
-\lambda & 1-\lambda^2/2 & A\lambda^2 \\
A\lambda^3 (1-\rho - i\eta) & -A\lambda^2 & 1
\end{pmatrix},$$

where $A$ and $|\rho + i\eta|$ are $O(1)$, while the mixing parameter $\lambda$ is assumed to be small. For the masses $m_i$ and the eigenvalues $M_i$ of the Majorana mass matrix $M$ we assume the same hierarchy which is observed for up-type quarks,

$$m_1 = b\lambda^4 m_3, \quad m_2 = c\lambda^2 m_3, \quad b, c = O(1)$$

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For the eigenvalues $m_i$ of the Dirac mass matrix this choice is suggested by SO(10) unification. For the masses $M_i$ this is an assumption motivated by simplicity. The masses $M_i$ cannot be degenerate, because in this case there exists a basis for $\nu_\alpha$ such that $U = 1$, which implies that no baryon asymmetry is generated. However, the precise form of the assumed hierarchy has no influence on the viability of the leptogenesis mechanism\textsuperscript{[5]}. The see-saw mechanism then yields light neutrino masses,

$$m_{\nu_\alpha} = \frac{b^2}{|C + e^{4i\alpha} B|} \lambda^4 m_{\nu_\alpha} + O(\lambda^6),$$

where the eigenvalues $m_i$ of $m_{D,\text{diag}}$ are real and positive. In the absence of a Majorana mass term $V$ and $U$ would correspond to Kobayashi-Maskawa type mixing matrices of left- and right-handed charged currents, respectively.

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where the eigenvalues $m_i$ of $m_{D,\text{diag}}$ are real and positive. In the absence of a Majorana mass term $V$ and $U$ would correspond to Kobayashi-Maskawa type mixing matrices of left- and right-handed charged currents, respectively.
\begin{align}
  m_{\nu}\mu & = \frac{c^2 |C + e^{i\alpha} B|}{BC} \lambda^2 m_{\nu}\tau + \mathcal{O}\left(\lambda^4\right), \\
  m_{\nu}\tau & = \frac{m_3^2}{M_3} + \mathcal{O}\left(\lambda^4\right). \tag{12}
\end{align}

The \(CP\)-asymmetry in the decay of the lightest right-handed neutrino \(N_1\) is easily obtained from eqs. (3) and (8)-(10),
\begin{align}
  \varepsilon_1 = -3 \frac{B A^2}{16\pi c^2 + A^2} \lambda^4 \frac{m_3^2}{v^2} \text{Im} \left[ (\rho - i\eta)^2 e^{i2(\alpha + \sqrt{3}\beta)} \right] + \mathcal{O}\left(\lambda^6\right). \tag{14}
\end{align}

For \(\lambda \sim 0.1\) one has \(|\varepsilon_1| \leq 10^{-6} \cdot m_3^2/v^2\). Hence, a large value of the Yukawa coupling \(m_3/v\) will be required by this mechanism of baryogenesis. This holds irrespective of the neutrino mixings and the heavy neutrino masses.

3. Numerical Results

To obtain a numerical value for the produced baryon asymmetry, one has to specify the free parameters in the ansatz (8)-(10). In the following we will use as a constraint the value for the \(\nu_\mu\)-mass which is preferred by the MSW explanation\(^{10}\) of the solar neutrino deficit,
\begin{align}
  m_{\nu}\mu & \simeq 3 \cdot 10^{-3} \text{ eV}. \tag{15}
\end{align}

A generic choice for the free parameters is to take all \(\mathcal{O}(1)\) parameters equal to one and to fix \(\lambda\) to a value which is of the same order as the \(\lambda\) parameter of the quark mixing matrix,
\begin{align}
  A = B = C = b = c = |\rho + i\eta| \simeq 1, \quad \lambda \simeq 0.1. \tag{16}
\end{align}

From eqs. (11)-(13), (13) and (16) one now obtains,
\begin{align}
  m_{\nu}\mu & \simeq 8 \cdot 10^{-6} \text{ eV}, \quad m_{\nu}\tau \simeq 0.15 \text{ eV}. \tag{17}
\end{align}

Finally, a second mass scale has to be specified. In unified theories based on SO(10) the Dirac neutrino mass \(m_3\) is naturally equal to the top-quark mass,
\begin{align}
  m_3 = m_t \simeq 174 \text{ GeV}. \tag{18}
\end{align}

This determines the masses of the heavy Majorana neutrinos \(N_i\), \(M_1 \simeq 2 \cdot 10^{14} \text{ GeV}\) and, consequently, \(M_1 \simeq 2 \cdot 10^{12} \text{ GeV}\) and \(M_2 \simeq 2 \cdot 10^{12} \text{ GeV}\). From eq. (14) one obtains the \(CP\) asymmetry \(|\varepsilon_1| \simeq 3 \cdot 10^{-6}\), where we have assumed maximal phases. The solution of the Boltzmann equations now yields the \((B - L)\) asymmetry,
\begin{align}
  Y_{B-L} \simeq 8 \cdot 10^{-10}, \tag{19}
\end{align}

which is indeed the correct order of magnitude!
Figure 4: Decay modes of the right-handed Neutrino superfield.

The large mass $M_3$ of the heavy Majorana neutrino $N_3$ suggests that $B - L$ is already broken at the unification scale $\Lambda_{\text{GUT}} \sim 10^{16}$ GeV, without any intermediate scale of symmetry breaking. This large value of $M_3$ is a consequence of the choice $m_3 \simeq m_t$. To test the sensitivity of the result for $Y_{B-L}$ on this assumption, consider as an alternative the choice $m_3 = m_b \simeq 4.5$ GeV, with all other parameters remaining unchanged. In this case one has $M_3 = 10^{11}$ GeV and $|\epsilon_1| = 2 \cdot 10^{-9}$ for the mass of $N_3$ and the $CP$ asymmetry, respectively. Since the maximal $B - L$ asymmetry is $-\epsilon_1/g^*$ (cf. [1]), it is clear why the generated asymmetry,

$$Y_{B-L} \simeq 6 \cdot 10^{-13},$$

is too small by more than two orders of magnitude. We conclude that high values for both masses $m_3$ and $M_3$ are preferred, which is natural in SO(10) unification.

Models for dark matter involving massive neutrinos favour a $\tau$-neutrino mass $m_\nu_\tau \simeq 5 \text{ eV}$ [2], which is significantly larger than the value given in (17). Such a large value for the $\tau$-neutrino mass can be accommodated within the ansatz described in this section. However, it does not correspond to the simplest choice of parameters and requires some fine-tuning. For the mass of the heaviest Majorana neutrino one obtains in this case $M_3 \simeq 6 \cdot 10^{12}$ GeV.

4. Supersymmetric extension

Without an intermediate scale of symmetry breaking, the unification of gauge couplings appears to require low-energy supersymmetry. Supersymmetric leptogenesis has already been considered in an approximation where lepton number violating scatterings are neglected which inhibit the generation of lepton number. However, a full analysis of the mechanism including all the relevant scattering processes is necessary in order to get a reliable relation between the input parameters and the final asymmetry. It turns out that the lepton number violating scatterings are qualitatively more important than in the non-supersymmetric scenario and that they can even account for the generation of an equilibrium distribution of heavy neutrinos at high temperatures.
The supersymmetric generalization of the Lagrangian (1) is the superpotential

$$W = \frac{1}{2} N^c M N^c + \mu H_1 \epsilon H_2 + H_1 \epsilon L \lambda_1 E^c + H_2 \epsilon L \lambda_2 N^c,$$

(21)

where, in the usual notation, $H_1$, $H_2$, $L$, $E^c$ and $N^c$ are chiral superfields describing spin-0 and spin-$\frac{1}{2}$ fields. The basis for the lepton fields can be chosen as in the non-supersymmetric case. The vacuum expectation values $v_1 = \langle H_1 \rangle$ and $v_2 = \langle H_2 \rangle$ of the two neutral Higgs fields generate Dirac masses $m_l = \lambda_1 v_1$ for the charged leptons and their scalar partners, and $m_\nu = \lambda_2 v_2$ for the neutrinos. The heavy neutrinos and their scalar partners can decay into various final states which can be summarized in the superfield diagrams in fig. 4. At tree level, the decay widths read,

$$\Gamma_{rs} (N_i \rightarrow \tilde{\ell} + \tilde{h}^c) = \Gamma_{rs} (N_i \rightarrow l + H_2) = \frac{(m^D_D^\dagger m_D)_ii}{16\pi v^2} M_i,$$

(22)

$$\Gamma_{rs} (\tilde{N}_i \rightarrow \tilde{l} + H_2) = \Gamma_{rs} (\tilde{N}_i \rightarrow \tilde{l} + \tilde{h}^c) = \frac{(m^D_D^\dagger m_D)_ii}{8\pi v^2} M_i.$$

(23)

The $\text{CP}$ asymmetry in each of the decay channels is given by

$$\varepsilon_i = \frac{1}{8\pi v^2} \sum_{j \neq i} \text{Im} \left[ \frac{(m^D_D^\dagger m_D)_ij^2}{(m^D_D^\dagger m_D)_ii} \right] g \left( \frac{M_j^2}{M_i^2} \right).$$

(24)

$$g(x) = \sqrt{x} \left[ \ln \left( \frac{1+x}{x} \right) + \frac{2}{x-1} \right].$$

(25)

It arises through interference of tree level and one-loop diagrams shown in fig. 4. In the case of a mass hierarchy, $M_j \gg M_i$, the $\text{CP}$ asymmetry is twice as large as in the non-supersymmetric case.

Like in the non-supersymmetric scenario lepton number violating scatterings mediated by a heavy $(s)$neutrino have to be included in a consistent analysis, since they can easily reduce the generated asymmetry by two orders of magnitude. A very interesting new feature of the supersymmetric model is that the $(s)$neutrino $(s)$top scatterings are strong enough to bring the neutrinos into thermal equilibrium at high temperatures. Hence, an equilibrium distribution can be reached for temperatures far below the masses of heavy gauge bosons.

Further, it turns out that the asymmetry essentially depends on the ratio

$$\tilde{m}_1 = \frac{(m^D_D^\dagger m_D)_ii}{M_1}.$$

(26)

For the mass matrices discussed in sect. 4 $\tilde{m}_1$ is of the same order as the muon neutrino mass. One easily verifies,

$$\tilde{m}_1 = \frac{C(c^2 + A^2 |\rho + i\eta|^2)}{c^2 |C + e^{4\alpha} B|} m_{\nu_\mu} + \mathcal{O}(\lambda^2).$$

(27)
In fig. 5 we have plotted the generated lepton asymmetry as function of $\tilde{m}_1$ for three different values of $M_1$, where we have assumed the hierarchy $M_2 = 10^3 M_1$, $M_3 = 10^6 M_1$ and the $CP$ asymmetry $\varepsilon_1 = -10^{-6}$.

Fig. 5 demonstrates, that in the whole parameter range the generated asymmetry is much smaller than the value $\varepsilon/g_* \sim 4 \cdot 10^{-9}$ which one obtains by neglecting lepton number violating scattering processes. For small $\tilde{m}_1$ the reason is that the Yukawa interactions are too weak to bring the neutrinos into equilibrium at high temperatures. For large $\tilde{m}_1$, on the other hand, the lepton number violating scatterings wash out a large part of the generated asymmetry at temperatures $T < M_1$.

Baryogenesis is possible in the range

$$10^{-5} \text{ eV} \lesssim \tilde{m}_1 \lesssim 5 \cdot 10^{-3} \text{ eV}.$$  \hspace{1cm} (28)

This result is independent of any assumptions on the mass matrices, in particular it is not a consequence of the ansatz discussed in sect. 5. This ansatz just implies $\tilde{m}_1 \approx m_{\nu_\mu}$ (cf. (27)). It is very interesting that the $\nu_\mu$-mass preferred by the MSW explanation of the solar neutrino deficit lies indeed in the interval allowed by baryogenesis according to fig. 5.

Consider now again the simplest choice of parameters given by eqs. (15)-(18). The corresponding generated lepton asymmetries are shown in fig. 6a. $Y_{L f}$ and $Y_{L s}$ denote the absolute values of the asymmetries stored in leptons and their scalar partners,
Figure 6: Generated asymmetry if one assumes a similar pattern of masses and mixings for the leptons and the quarks. In both figures we have $\lambda = 0.1$ and $m_3 = m_t$ (a) and $m_3 = m_b$ (b). The hatched area shows the lepton asymmetry corresponding to the measured baryon asymmetry.

respectively. They are related to the baryon asymmetry by

$$Y_B = -\frac{8}{23} Y_L, \quad Y_L = Y_{L_f} + Y_{L_s}. \quad (29)$$

$Y_{N_1}$ is the number of heavy neutrinos per comoving volume element, and $Y_{1 \pm} = Y_{N_1} \mp Y_{N_1} \equiv Y_{\tilde{N}_1}$, where $Y_{\tilde{N}_1}$ is the number of scalar neutrinos per comoving volume element. As fig. (6a) shows, the generated baryon asymmetry has the correct order of magnitude,

$$Y_{B-L} \simeq 1 \cdot 10^{-9}. \quad (30)$$

Lowering the Dirac mass scale of the neutrinos to the bottom-quark scale has again dramatic consequences. The baryon asymmetry is reduced by three orders of magnitude

$$Y_{B-L} \simeq 1 \cdot 10^{-12}. \quad (31)$$

Hence, like in the non-supersymmetric scenario, large values for both masses $m_3$ and $M_3$ are necessary.

Comparing the results (30) and (31) with their non-supersymmetric counterparts (19) and (20), one sees that the larger $CP$ asymmetry and the additional contributions from the sneutrino decays in the supersymmetric scenario are compensated by the wash-out processes which are stronger than in the non-supersymmetric case. The final asymmetries are the same in the non-supersymmetric and in the supersymmetric case.
The recently reported atmospheric neutrino anomaly may be due to neutrino oscillations. The required mass difference and mixing angle are $\Delta m^2 \sim 0.005 \text{ eV}^2$ and $\sin^2 2\Theta \sim 1$. The preferred solution for baryogenesis discussed above yields (cf. eq. (17)) $m_{\nu_\tau}^2 - m_{\nu_\mu}^2 \simeq 0.02 \text{ eV}^2$ which, within the theoretical and experimental uncertainties, is certainly consistent with the mass difference required by the neutrino oscillation hypothesis. The $\nu_\tau$-$\nu_\mu$ mixing angle is not constrained by leptogenesis and therefore a free parameter in principle. The large value needed, however, is against the spirit of small generation mixings manifest in the Wolfenstein ansatz and would require some special justification.

5. Summary

Anomalous electroweak $B + L$ violating processes are in thermal equilibrium in the high-temperature phase of the standard model. Hence, the cosmological baryon asymmetry can be generated from a primordial lepton asymmetry. Necessary ingredients are right-handed neutrinos and Majorana masses, which occur naturally in SO(10) unification.

The baryon asymmetry can be computed by standard methods based on Boltzmann equations. In a consistent analysis lepton number violating scatterings have to be taken into account, since they can erase a large part of the asymmetry. In supersymmetric scenarios these scatterings are sufficient to generate an initial equilibrium distribution of heavy Majorana neutrinos.

Baryogenesis implies stringent constraints on the light neutrino masses. Assuming a similar pattern of mixings and masses for neutrinos and up-type quarks, as suggested by SO(10) unification, the observed asymmetry is obtained without any fine tuning. The $\nu_\mu$ mass is predicted in a range consistent with the MSW solution of the solar neutrino problem. $B - L$ is broken at the unification scale. The baryogenesis scale is given by the mass of the lightest of the heavy Majorana neutrinos, which is much lower and consistent with constraints from inflation and the gravitino abundance.

As our discussion illustrates, the cosmological baryon asymmetry is closely related to neutrino properties. Already the existence of a baryon asymmetry is a strong argument for lepton number violation and Majorana neutrino masses. Together with further information about neutrino properties from high-energy physics and astrophysics, the theory of the baryon asymmetry will give us new insights into physics beyond the standard model.
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