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Baryogenesis and the scale of $B-L$ breaking

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In unified theories with right-handed neutrinos the scale of $B-L$ breaking can a priori vary between the unification scale and the Fermi scale of electroweak symmetry breaking. We derive upper and lower bounds on the masses of light and heavy Majorana neutrinos and on the scale of $B-L$ breaking from the requirement that the cosmological baryon asymmetry is generated before the electroweak phase transition. Baryogenesis poses a serious problem for extended gauge theories with heavy neutrinos and neutral vector bosons in the TeV mass range.

In the standard model of strong and electroweak interactions baryon number B and the three lepton numbers L_e, L_μ, and L_τ correspond to global symmetries of the classical action. However, $B+L$ and $B-L$, where $L=L_e+L_\mu+L_\tau$ is the total lepton number, are broken by gauge and gravitational anomalies, and only L_e-L_μ and $L_\mu-L_\tau$ are conserved charges of the quantum theory. On the contrary, in unified theories based on the gauge group SO(10) or E_6, which contain right-handed neutrinos in addition to the quarks and leptons of the standard model, $B-L$ plays the role of a spontaneously broken local symmetry, and neutrinos acquire masses related to the scale of $B-L$ breaking. There are no remaining exact global symmetries. This is a theoretically very appealing extension of the standard model which implies lepton number violating processes at low and high energies. The mass scale of $B-L$ breaking can a priori vary between the unification scale and the Fermi scale of electroweak symmetry breaking.

In the very early universe, at temperatures above the critical temperature of the electroweak phase transition, $B+L$ violating interactions were in thermal equilibrium [1]. As a consequence, the equilibrium values of B and $B-L$ were proportional,

$$\langle B\rangle_T=\gamma\langle B-L\rangle_T\,, \tag{1}$$

where the constant γ depends on the particle content of the theory [2,3]. Hence, the generation of a cosmological baryon asymmetry, an important postdiction of grand unifield theories, was strongly dependent on the strength of $B-L$ violating interactions [3–5] and, in particular, on the mass scale of spontaneous $B-L$ breaking.

According to eq. (1), generation of a baryon asymmetry requires the breaking of $B-L$. Recently, extended gauge theories with a mass scale Λ_{B-L} in the TeV range have received much attention, mostly in connection with superstring theories. The corresponding vector bosons and fermions with masses between a few tens of GeV and a few TeV can be searched for at present and future e^+e^-, ep and pp colliders [6]. In these theories the observed baryon asymmetry could in principle be produced via the out-of-equilibrium decay of heavy neutrinos [7]. The generated lepton asymmetry then gives rise to a baryon asymmetry due to $B+L$ violating processes which are in thermal equilibrium above the critical temperature of the electroweak phase transition. This mechanism has recently been studied in detail [8]. There are also suggestions to generate a baryon asymmetry during the electroweak phase transition [9]. However, the proposed mechanisms all depend on details of the phase transition, and at present it is unclear whether they can yield the required amount of baryon asymmetry.

In the following we shall consider the constraints on light and heavy neutrino masses and on the scale of $B-L$ breaking, which have to be satisfied it the baryon asymmetry is generated by out-of-equilibrium decays of heavy neutrinos. The leptonic part of the theory is described by the lagrangian

$$\mathscr{L}_l = \bar{l}\,\mathrm{i}\not{D}l + \bar{\nu}_\mathrm{R}\,\mathrm{i}\not{\partial}\nu_\mathrm{R} - \bar{l}\phi g_\nu \nu_\mathrm{R} - \bar{\nu}_\mathrm{R} g_\nu^\dagger \phi^\dagger l - \tfrac{1}{2}(\bar{\nu}_\mathrm{R} M \nu_\mathrm{R}^\mathrm{c} + \bar{\nu}_\mathrm{R}^\mathrm{c} M^\dagger \nu_\mathrm{R}) \,. \tag{2}$$

Here $l=(\nu_\mathrm{L}, e_\mathrm{L}^-)$ is the lepton doublet, ϕ is the doublet of Higgs fields, D_μ is the gauge covariant derivative and $\nu_\mathrm{R}^\mathrm{c} = C\bar{\nu}_\mathrm{R}^\mathrm{T}$, where C is the charge conjugation matrix. The mass matrix M is generated by spontaneous symmetry breaking, i.e., $M=hv'$, where h is a matrix of Yukawa couplings and v' is the vacuum expectation value which breaks the extended gauge symmetry to the standard model gauge group. For simplicity, we shall restrict ourselves to the case of two families, i.e., g_ν and M are 2×2 complex matrices. One can always choose a basis for ν_R such that

$$M = \begin{pmatrix} M_1 & 0 \\ 0 & M_2 \end{pmatrix}, \tag{3}$$

where M_1 and M_2 are real and positive. The Dirac mass matrix, which is generated after the spontaneous breaking of the electroweak gauge group, has the form

$$m_\mathrm{D} = g_\nu v = V\tilde{m}_\mathrm{D} \,, \tag{4}$$

where V is the Kobayashi–Maskawa type matrix in the leptonic charged current and $v=174$ GeV. The matrix $\tilde{m}_\mathrm{D}$ enters the see-saw formula [10] for the light neutrino masses which, for two generations, takes the form

$$\begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix} = \tilde{m}_\mathrm{D} \frac{1}{M} \tilde{m}_\mathrm{D}^\mathrm{T} \tag{5}$$

$$= \begin{pmatrix} a^2/M_1 & 0 \\ 0 & d^2/M_2 \end{pmatrix} (1+\eta^2) \,, \tag{6}$$

where

$$\tilde{m}_\mathrm{D} = \begin{pmatrix} a & b \\ c & d \end{pmatrix}, \tag{7}$$

$$b = \sqrt{\frac{M_2}{M_1}}\,\eta a, \quad c = \sqrt{\frac{M_1}{M_2}}\,\eta d \,, \tag{8}$$

$$(1+\eta^2) = \zeta \exp(\mathrm{i}\delta) \,, \tag{9}$$

ζ can be any positive real number. The matrix elements of $\tilde{m}_\mathrm{D}$ are not all independent since the matrix (5) has to be diagonal. Note, that for $\zeta \ll 1$ one can have very small light neutrino masses but nevertheless large mixings between light and heavy neutrinos in charged and neutral currents, which are given by the matrix $\xi = \tilde{m}_\mathrm{D}/M$ [11]. For $M_1 \sim 100$ GeV, a mixing angle $\xi_{ij}^2 \sim 10^{-2}$ requires typically $\zeta \sim 10^{-6}$, such that the light neutrino masses satisfy the upper experimental bounds $m_{\nu_e} < 8$ eV, $m_{\nu_\mu} < 270$ keV [12].

Let us now consider the decay of the heavy neutrino N_1 (cf. fig. 1). The out-of-equilibrium condition is satisfied if the decay rate Γ_1 of N_1 is smaller than the Hubble parameter at temperature $T=M_1$ (cf. ref. [13]),

$$\Gamma_1/2H(T=M_1) < 1 \,, \tag{10}$$

the Hubble parameter $H=1.7\, g_*^{1/2} T^2/m_\mathrm{Pl}$, where $g_* \approx 100$ is the effective number of degrees of freedom of the standard model and $m_\mathrm{Pl} = 1.2\times10^{19}$ GeV. The decay rate Γ_1 is given by

$$\Gamma_1 = \frac{1}{16\pi} (g_\nu^\dagger g_\nu)_{11} M_1 \,. \tag{11}$$

Using eqs. (4)–(11) one obtains the following constraint on the neutrino masses:

$$\frac{m_1 + m_2|\eta|^2}{|1+\eta^2|} = \frac{1}{M_1} (\tilde{m}_\mathrm{D}^\dagger m_\mathrm{D})_{11} < 4\times10^{-3}\ \mathrm{eV} \,. \tag{12}$$

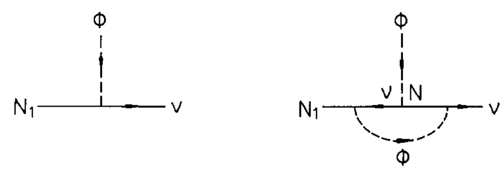

Fig. 1. Decay of a heavy Majorana neutrino into a massless lepton and a massless scalar particle.

Furthermore, N_1 has to decay sufficiently fast, i.e., before the temperature T_{SPH} is reached above which $B+L$ changing interactions are in thermal equilibrium. This means

$$M_1 > T_{SPH} , \tag{13}$$

where the temperature T_{SPH} is of the order of the zero-temperature vacuum expectation value v.

A third condition on the neutrino masses follows from the requirement that the lepton asymmetry, which is generated by the decay of the heavy neutrino N_1 and converted into a baryon asymmetry via sphaleron processes, is large enough. The strength of the CP-violating interference term of tree-level and one-loop amplitudes (cf. fig. 1) is given by [7]

$$\epsilon = \frac{\mathrm{Im}[(\tilde{m}_D^\dagger \tilde{m}_D)_{12}^2 I(M_2/M_1)]}{\pi v^2 (\tilde{m}_D^\dagger \tilde{m}_D)_{11}} , \tag{14}$$

where

$$I(x) = x\left(1 + (1+x^2)\ln\frac{x^2}{1+x^2}\right) \tag{15}$$

$$\approx 1/2x, \quad x \gg 1 . \tag{16}$$

The corresponding baryon asymmetry is (cf. ref. [13])

$$B \equiv \frac{n_B}{g_* n_\gamma} \approx \frac{\epsilon}{g_*} , \tag{17}$$

where n_B and n_γ are baryon and photon number densities, respectively. For $\epsilon > \epsilon_0 \approx 10^{-8}$ a large enough baryon asymmetry can be generated. From eqs. (4)–(9) and (14)–(16) one obtains

$$\epsilon = \frac{\sin\delta M_1 (m_1+m_2)^2}{2\pi v^2 (m_1 + m_2 |\eta|^2)} > \epsilon_0 . \tag{18}$$

Finally, neutrino masses are restricted by the requirement that $\Delta L=2$ processes (cf. fig. 2) do not erase a generated lepton asymmetry [3,4]. From the out-of-equilibrium condition $\Gamma_{\Delta L=2}/2H(T=M_1) \ll 1$ and $\Gamma_{\Delta L=2} = \sum m_\nu^2 T^3/\pi^3 v^4$ one obtains the constraint

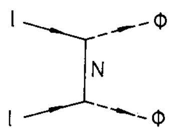

Fig. 2. $\Delta L=2$ scattering process.

$$m_1^2 + m_2^2 < (30\,\mathrm{keV})^2 \left(\frac{100\,\mathrm{GeV}}{M_1}\right) , \tag{19}$$

where we have set the temperature at which the lepton asymmetry is generated equal to M_1.

Let us now consider the implications of the constraints (12), (18) and (19) on light and heavy Majorana neutrino masses. $m_-(m_+)$ be the smaller (larger) mass of the two masses m_1 and m_2. The out-of-equilibrium condition (12) immediately yields

$$m_- < 4\times 10^{-3}\,\mathrm{eV} . \tag{20}$$

This agrees with the bound obtained by Fischler et al. [4] in the case of one generation. It applies for the lightest neutrino in the case of two generations. Eq. (20) is an important prediction of the considered mechanism of baryogenesis which is independent of the scale of $B-L$ breaking.

The lower bound on the heavy neutrino masses is strongly dependent on the structure of $\tilde{m}_D$. In the case of large mixing between light and heavy neutrinos in neutral and charged currents one has $\zeta \ll 1$ (cf. eq. (9)) and therefore $|\eta|^2 \approx 1$. In this case the out-of-equilibrium condition (12) and the baryogenesis condition (18) imply the stringent bound

$$M_1 > 5\times 10^8\,\mathrm{GeV}\cdot\left(\frac{\epsilon}{\epsilon_0}\right)\frac{1}{\zeta} . \tag{21}$$

Hence, heavy neutrinos with masses of a few hundred GeV and large mixings with light neutrinos appear to be incompatible with baryogenesis! For $\zeta \ll 1$ one also obtains a stringent upper bound on the light neutrino masses from eq. (12),

$$m_+ < 4\times 10^{-3}\,\mathrm{eV}\cdot\zeta , \tag{22}$$

which is much more restrictive than the bound (20).

In the case of small mixings between light and heavy neutrinos ζ may be O(1) or larger, and a low value of M_1 requires a small mass ratio m_-/m_+. In order to derive the corresponding bound it is convenient to replace in eq. (18) $\sin\delta$ by $\sin\alpha\ |\eta|^2/\zeta$, where $\eta^2 = |\eta|^2 \exp(i\alpha)$. One then obtains

$$(m_1+m_2)^2 > 20\,\mathrm{keV}\left(\frac{100\,\mathrm{GeV}}{M_1}\right)\left(\frac{\epsilon}{\epsilon_0}\right) \times (m_1 + m_2|\eta|^2)\,\frac{|1+\eta^2|}{|\eta|^2} . \tag{23}$$

The two extreme cases are $m_1 \ll m_2 \equiv m_+$, for $|\eta|^2 \ll 1$, and $m_2 \ll m_1 \equiv m_+$ for $|\eta|^2 \gg 1$ (cf. eq. (12)). In both cases one finds

$$m_+ > 20\ \mathrm{keV}\left(\frac{100\ \mathrm{GeV}}{M_1}\right)\left(\frac{\epsilon}{\epsilon_0}\right). \quad (24)$$

For small values of M_1 eq. (24) is barely compatible with the condition (19) which guarantees that the generated lepton number is not erased. The bound (24) implies that the heavy neutrinos can only be within the reach of present and projected colliders if there is a very large hierarchy between the small neutrino masses ($m_-/m_+ < 10^{-7}$).

The upper bound on the cosmological mass density requires $m_+ < 100$ eV [13], unless one introduces new particles which allow a sufficiently fast decay of ν_+. Together with eq. (24) this implies

$$M_1 > 2\times 10^4\ \mathrm{GeV}\cdot\left(\frac{\epsilon}{\epsilon_0}\right). \quad (25)$$

A much more stringent bound obtains if the MSW mechanism is responsible for the observed solar neutrino flux. Identifying m_+^2 with the fit of the GALLEX results $\Delta m^2 \approx 7\times 10^{-6}\ \mathrm{eV}^2$ [14] one finds, using eq. (24),

$$M_1 > 7\times 10^8\ \mathrm{GeV}\cdot\left(\frac{\epsilon}{\epsilon_0}\right), \quad (26)$$

which is comparable to the bound (21).

So far, we have ignored possible new gauge interactions of the heavy neutrinos which are present if $B-L$ is a spontaneously broken, local symmetry. There could be additional neutral Z' bosons or, in left–right symmetric models, also charged W_R bosons. As recently shown in ref. [15], these gauge bosons must be sufficiently heavy in order to have enough departure from thermal equilibrium such that a lepton asymmetry can be generated. A simple estimate for the allowed range of gauge boson masses M_G is obtained from the out-of-equilibrium condition $\Gamma_G/2H(T=M_1) < 1$. The rate for gauge boson mediated processes, such as $Ne_R \to d_R u_R^c$ or $NN \to ee^c$, is $\Gamma_G \sim g^4 M_1^5/4\pi M_G^4$. With $g \approx 0.5$, one finds the lower bound on M_G (cf. ref. [8]):

$$M_G > 10^5\ \mathrm{GeV}\left(\frac{M_1}{100\ \mathrm{GeV}}\right)^{3/4}. \quad (27)$$

Hence, in unified gauge theories, where $B-L$ is a local symmetry, the scale of $B-L$ breaking must always be much larger than the Fermi scale of electroweak symmetry breaking. If the effect of new gauge interactions is ignored one reaches a different conclusion [16].

More precise bounds on light and heavy neutrino masses can be obtained from an integration of the Boltzmann equations. In particular the out-of-equilibrium condition is then effectively replaced by $\Gamma_1/2H(T=M_1) = K < O(100)$ [13,8]. A simple analysis shows that the upper bounds on the light neutrino masses are roughly multiplied by K, whereas the lower bounds on the heavy neutrino masses remain essentially unchanged.

To summarize. In extended gauge theories, where $B-L$ is spontaneously broken at a mass scale Λ_{B-L} below the unification scale Λ_{GUT}, the cosmological baryon asymmetry has to be generated at temperatures below Λ_{B-L}. At present, the only mechanism for which it has been demonstrated that a large enough baryon asymmetry can be produced is the decay of heavy neutrinos. Our analysis shows that this mechanism requires heavy neutrinos with masses far above the Fermi scale, unless one allows for extreme fine tuning of parameters in the neutrino mass matrix. Furthermore, the masses of new vector bosons, and therefore the mass scale of $B-L$ breaking, must be much larger than the smallest heavy neutrino mass. Hence, baryogenesis poses a serious problem for extended gauge theories with heavy neutrinos and new neutral vector bosons whose masses are in the TeV range. The discovery of such particles at present and future colliders would strongly suggest a novel mechanism of baryogenesis, possible related to the electroweak phase transition.

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