BARYOGENESIS ABOVE THE FERMI SCALE

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Abstract
In the standard model and most of its extensions the electroweak transition is too weak to affect the cosmological baryon asymmetry. Due to sphaleron processes baryogenesis in the high-temperature, symmetric phase of the standard model is closely related to neutrino properties. The experimental indications for very small neutrino masses from the solar and the atmospheric neutrino deficits favour a large scale of $B - L$ breaking. For hierarchical neutrino masses, with $B - L$ broken at the unification scale $\Lambda_{\text{GUT}} \sim 10^{16}$ GeV, the observed baryon asymmetry $n_B/s \sim 10^{-10}$ is naturally explained by the decay of heavy Majorana neutrinos. The corresponding baryogenesis temperature is $T_B \sim 10^{10}$ GeV. In supersymmetric models implications for the mass spectrum of superparticles can be derived. A consistent picture is obtained with the gravitino as LSP, which may be the dominant component of cold dark matter.

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1. General aspects of baryogenesis

The cosmological baryon asymmetry, the ratio of the baryon density to the entropy density of the universe,
\[ Y_B = \frac{n_B}{s} = (0.6 - 1) \cdot 10^{-10}, \tag{1} \]
can be understood as a consequence of baryon number violation, C and CP violation, and a deviation from thermal equilibrium. 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. 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. 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 (cf. fig. 1) which violates baryon and lepton number by three units,
\[ \Delta B = \Delta L = 3. \tag{2} \]
The evaluation of the sphaleron rate in the symmetric high temperature phase is a challenging problem. Although a complete theoretical understanding has not yet been achieved, it is generally believed that \( B \) and \( L \) violating processes are in thermal equilibrium for temperatures in the range
\[ T_{EW}^c \sim 100 \text{ GeV} < T < T_{SPH} \sim 10^{12} \text{ GeV}. \tag{3} \]
Sphaleron processes have 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 \( T > T_{EW}^c \), 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 \( Y_B = (n_B - \overline{n_B})/s \) and the corresponding \( L \) and \( B - L \) asymmetries \( Y_L \) and \( Y_{B-L} \), respectively,
\[ Y_B = C Y_{B-L} = \frac{C}{C - 1} Y_L. \tag{4} \]
Here \( C \) is a number \( \mathcal{O}(1) \). In the standard model with three generations and one Higgs doublet one has \( C = 28/79 \).
We conclude that $B - L$ violation is needed if the baryon asymmetry is generated before the electroweak transition, i.e. at temperatures $T > T_{EW}^c \sim 100$ 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 in extensions of the standard model at low energies is the appearance of an effective $\Delta L = 2$ interaction between lepton and Higgs fields,

$$\mathcal{L}_{\Delta L=2} = \frac{1}{2} \overline{l_L} \phi g_\nu \frac{1}{M} g^T_\nu \phi l_L + \text{h.c.} .$$

Such an interaction arises in particular from the exchange of heavy Majorana neutrinos (cf. fig. 2). 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 $\nu_\tau$.

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

**Figure 2:** *Lepton number violating lepton Higgs scattering*
rate

\[ \Gamma_{\Delta L=2}(T) = \frac{1}{\pi^3} \frac{T^3}{v^4} \sum_{i=e,\mu,\tau} m_{\nu_i}^2. \]  

(6)

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

\[ Y_B = Y_{B-L} = Y_L = 0. \]  

(7)

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

\[ \sum_{i=e,\mu,\tau} m_{\nu_i}^2 < (0.2 \text{ eV})^2. \]  

(8)

This bound is comparable to the upper bound on the electron neutrino mass obtained from neutrinoless double beta decay. Note, however, that the bound also applies to the \( \tau \)-neutrino mass. In supersymmetric theories two chiral U(1) symmetries in addition to baryon and lepton number are approximately conserved at temperatures above \( T_{SS} \sim 10^7 \text{ GeV} \). This relaxes the upper bound (8) from 0.2 eV to about 60 eV.

The connection between lepton number and the baryon asymmetry is lost if baryogenesis takes place at or below the Fermi scale. 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. In the minimal supersymmetric extension of the standard model (MSSM) such a scenario appears still possible for a limited range of parameters.

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 theories based on gauge groups containing the Pati-Salam symmetry, 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 reads

\[ \mathcal{L}_Y = -\overline{l_L} \tilde{\phi} g_l e_R - \overline{l_L} \phi g_e \nu_R - \frac{1}{2} \nu_R^\dagger M \nu_R + \text{h.c.}. \]  

(9)

The vacuum expectation value of the Higgs field \( \langle \varphi^0 \rangle = v \neq 0 \) generates Dirac masses \( m_l \) and \( m_D \) for charged leptons and neutrinos, \( m_l = g_l v \) and \( m_D = g_e v \), respectively,
which are assumed to be much smaller than the Majorana masses $M$. This yields light and heavy neutrino mass eigenstates

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

(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. The decay width of the heavy neutrino $N_i$ reads at tree level,

$$\Gamma_{Di} = \Gamma \left( N_i \to \phi^c + l \right) + \Gamma \left( N_i \to \phi + l^c \right)
= \frac{1}{8\pi} \frac{(m_D^\dagger m_D)_{ii}}{v^2} M_i.$$  

(12)

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

$$\tilde{m}_1 = \frac{(m_D^\dagger m_D)_{11}}{M_1} < 10^{-3} \text{eV}.$$  

(13)

More direct bounds on the light neutrino masses depend on the structure of the Dirac neutrino mass matrix as we shall discuss below.

Interference between the tree-level amplitude and the one-loop self-energy and vertex corrections yields the $CP$ asymmetry,

$$\varepsilon_1 = \frac{\Gamma(N_1 \to l\phi^c) - \Gamma(N_1 \to l^c\phi)}{\Gamma(N_1 \to l\phi^c) + \Gamma(N_1 \to l^c\phi)}
\simeq \frac{3}{16\pi v^2} \left. \frac{1}{(m_D^\dagger m_D)_{11}} \sum_{i=2,3} \text{Im} \left[ (m_D^\dagger m_D)_{ii}^2 \right] \frac{M_i}{M_1} \right\}.  

(14)

Here we have assumed $M_1 \ll M_2, M_3$. For very small mass differences, which are comparable to the decay widths, one obtains a resonance enhancement.

The $CP$ asymmetry (14) leads to the generated lepton asymmetry,

$$Y_L = \frac{n_L - n_R}{s} = \kappa \frac{\varepsilon_1}{g^*}.$$  

(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. In the examples discussed below one has $\kappa \simeq 0.1 \ldots 0.01.$
The CP asymmetry (14) is given in terms of the Dirac and the Majorana neutrino mass matrices. One can always choose a basis for the right-handed neutrinos where the Majorana mass $M$ is diagonal with real and positive eigenvalues. $m_D$ is a general complex matrix, which can be diagonalized by a biunitary transformation. One then has

$$m_D = V \begin{pmatrix} m_1 & 0 & 0 \\ 0 & m_2 & 0 \\ 0 & 0 & m_3 \end{pmatrix} U^\dagger, \quad M = \begin{pmatrix} M_1 & 0 & 0 \\ 0 & M_2 & 0 \\ 0 & 0 & M_3 \end{pmatrix},$$

where $V$ and $U$ are unitary matrices and the $m_i$ 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.

Note, that according to eqs. (12) and (14) 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 $K$ 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 $CP$ violation and generation mixing which are relevant at high energies and at low energies, respectively.

In many models the quark and lepton mass hierarchies and mixings are parametrised in terms of a common mixing parameter $\lambda \approx 0.1$. Assuming a hierarchy for the right-handed neutrino masses similar to the one satisfied by up-type
Figure 4: Time evolution of the neutrino and the scalar neutrino number densities, and of the lepton asymmetries for $\lambda \simeq 0.1$ and $m_3 \simeq m_t$.

quarks,

$$\frac{M_1}{M_2} \sim \frac{M_2}{M_3} \sim \lambda^2,$$

and a corresponding CP asymmetry

$$\varepsilon_1 \sim \frac{\lambda^4}{16\pi} \frac{m_3^2}{v^2} \sim 10^{-6} \frac{m_3^2}{v^2},$$

one obtains indeed the correct order of magnitude for the baryon asymmetry if one chooses $m_3 \simeq m_t \simeq 174$ GeV, as expected in theories with Pati-Salam symmetry. Using as a constraint the value for the $\nu_\mu$-mass which is preferred by the MSW explanation of the solar neutrino deficit, $m_{\nu_\mu} \simeq 3 \cdot 10^{-3}$ eV, the ansatz implies for the other light and the heavy neutrino masses

$$m_{\nu_e} \simeq 8 \cdot 10^{-6} \text{ eV}, \quad m_{\nu_\tau} \simeq 0.15 \text{ eV}, \quad M_3 \simeq 2 \cdot 10^{14} \text{ GeV}.$$  \hspace{1cm} (19)

Consequently, one has $M_1 \simeq 2 \cdot 10^{16}$ GeV and $M_2 \simeq 2 \cdot 10^{12}$ GeV. The solution of the Boltzmann equations then yields the baryon asymmetry (see fig. 3),

$$Y_B \simeq 9 \cdot 10^{-11},$$

which is indeed the correct order of magnitude. The precise value depends on unknown phases.
The large mass $M_3$ of the heavy Majorana neutrino $N_3$ (cf. (19)), 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 $m_3 \simeq m_t$. This is indeed necessary in order to obtain sufficiently large CP asymmetry.

The recently reported atmospheric neutrino anomaly may be due to $\nu_\mu - \nu_\tau$ oscillations. The required mass difference is $\Delta m_{\nu_\mu \nu_\tau}^2 \simeq (5 \cdot 10^{-4} - 6 \cdot 10^{-3})$ eV$^2$, together with a large mixing angle $\sin^2 2\Theta > 0.82$. In the case of hierarchical neutrinos this corresponds to a $\tau$-neutrino mass $m_\nu \sim (0.02 - 0.08)$ eV. Within the theoretical uncertainties this is consistent with the $\tau$-neutrino mass (19) obtained from baryogenesis. The $\nu_\tau - \nu_\mu$ mixing angle is not constrained by leptogenesis and therefore a free parameter in principle. The large value, however, is different from the mixing angles known in the quark sector and requires an explanation. An possible framework are $U(1)$ family symmetries. A large mixing angle can also naturally occur together with a mass hierarchy of light and heavy Majorana neutrinos.

Without an intermediate scale of symmetry breaking, the unification of gauge couplings appears to require low-energy supersymmetry. Supersymmetric leptogenesis has recently been studied in detail, taking all relevant scattering processes into account, which is necessary in order to get a reliable relation between the input parameters and the final baryon asymmetry. It turns out that the lepton number violating scatterings are qualitatively more important than in the non-supersymmetric case and that they can also account for the generation of an equilibrium distribution of heavy neutrinos at high temperatures.

The supersymmetric generalization of the lagrangian (9) 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 vacuum expectation value $v_2 = \langle H_2 \rangle$ of the Higgs field $H_2$ generates the Dirac mass matrix $m_D = \lambda_\nu v_2$ for the neutrinos and their scalar partners.

The heavy neutrinos $N_i$ and their scalar partners $\tilde{N}_i$ decay with different probabilities into final states with different lepton number. The generated lepton asymmetries are shown in fig. $[\tilde{24}]$. $Y_{L_L}$ and $Y_{L_s}$ denote the absolute values of the asymmetries stored in leptons and their scalar partners, respectively. They are related to the baryon asymmetry by

$$Y_B = -\frac{8}{23} Y_L, \quad Y_L = Y_{L_L} + Y_{L_s}.$$

(22)

$Y_{N_i}$ is the number of heavy neutrinos per comoving volume element, and

$$Y_{1\pm} = Y_{\tilde{N}_i} \pm Y_{\tilde{N}_i},$$

(23)
where \( Y_{N_1} \) is the number of scalar neutrinos per comoving volume element. As fig. 4 shows, the generated baryon asymmetry has the correct order of magnitude, as in the non-supersymmetric case.

From the discussion of the out-of-equilibrium condition we know that the generated baryon asymmetry is very sensitive to the decay width \( \Gamma_{D_1} \) of \( N_1 \), and therefore to \( (m_D^\dagger m_D)_{11} \). In fact, the asymmetry essentially depends on the effective neutrino mass \( \tilde{m}_1 \). For the case of hierarchical neutrino masses described above, one has

\[
\tilde{m}_1 = \frac{(m_D^\dagger m_D)_{11}}{M_1} \simeq m_{\nu_\mu} .
\]  

(24)

It turns out that a sufficiently large baryon asymmetry is generated in the range

\[ 10^{-5} \text{ eV} \lesssim \tilde{m}_1 \lesssim 5 \cdot 10^{-3} \text{ eV} , \]

(25)

which is consistent with the rough bound (23). This result is independent of any assumptions on the mass matrices. 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!

Comparing non-supersymmetric and supersymmetric leptogenesis 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.

Leptogenesis can also be considered in extended models which contain heavy SU(2)-triplet Higgs fields in addition to right-handed neutrinos. Decays of the heavy scalar bosons can in principle also contribute to the baryon asymmetry. However, since these Higgs particles carry gauge quantum numbers they are strongly coupled to the plasma and it is difficult to satisfy the out-of-equilibrium condition. The resulting large baryogenesis temperature is in conflict with the ‘gravitino constraint’

3. SUSY mass spectrum and dark matter

The out-of-equilibrium condition for the decay of the heavy Majorana neutrinos, the see-saw mechanism and the experimental evidence for small neutrino masses are all consistent and suggest rather large heavy neutrino masses and a correspondingly large baryogenesis temperature. Within the ansatz described in the previous section one obtains

\[ T_B \sim M_1 \sim 10^{10} \text{ GeV} . \]

(26)

Such a large baryogenesis temperature can only be avoided in the very special case of a strong resonant amplification of the CP violating decays.

In the particularly attractive supersymmetric version of leptogenesis one also has to consider the following two issues: the size of other possible contributions to the
baryon asymmetry and the consistency of the large baryogenesis temperature with the ‘gravitino constraint\textsuperscript{43}. A large asymmetry may potentially be generated by coherent oscillations of scalar fields which carry baryon and lepton number\textsuperscript{31}. However, it appears likely that the interactions of the right-handed neutrinos are sufficient to erase such large primordial baryon and lepton asymmetries\textsuperscript{32}.

The ‘gravitino constraint’ is particularly interesting since it is model independent to a large extent, and it therefore provides very interesting information about possible extensions of the standard model. The production of gravitinos ($\tilde{G}$) at high temperatures is dominated by two-body processes involving gluinos ($\tilde{g}$) (cf. fig. 5). On dimensional grounds the production rate has the form

$$\Gamma(T) \propto \frac{1}{M^2} T^3,$$

(27)

where $M = (8\pi G_N)^{-1/2} = 2.4 \cdot 10^{18}$ GeV is the Planck mass. Hence, the density of thermally produced gravitinos increases strongly with temperature.

The production cross section is enhanced by a factor $(m_{\tilde{g}}/m_{\tilde{G}})^2$ for light gravitinos\textsuperscript{33}. The thermally averaged cross section has recently been evaluated for arbitrary gravitino masses. The result reads\textsuperscript{34}

$$C(T) = \langle \sigma_{(L)} v_{\text{rel}} \rangle$$

$$= \frac{21g^2(T)}{32\pi\zeta^2(3)M^2} \left((N^2 - 1)C_A + 2n_f NC_F\right) \left(1 + \frac{m_{\tilde{g}}^2(T)}{3m_{\tilde{G}}^2}\right)$$

$$\left(\ln \frac{1}{g^2(T)} + \frac{5}{2} + 2\ln 2 - 2\gamma_E\right).$$

(28)

Here $C_A$ and $C_F$ are the usual colour factors for the group SU(N) and $2n_f$ is the number of colour-triplet chiral multiplets, i.e. $2n_f = 12$ in the MSSM. The logarithmic collinear singularity of the cross section has been regularized by a plasma mass $m \sim g(T)T$ for the gluon. The unknown constant part of the thermally averaged cross
section is expected to be of the same size as the term proportional to ln(1/g^2(T)). For QCD (N=3) one has

\[ C(T) \simeq 10 \frac{g^2(T)}{M^2} \left( 1 + \frac{m_{\tilde{g}}^2(T)}{3m_{\tilde{G}}^2} \right) \left( \ln \frac{1}{g^2(T)} + 2.7 \right). \tag{29} \]

The cross section \( C(T) \) enters in the Boltzmann equation, which describes the generation of a gravitino density \( n_{\tilde{G}} \) in the thermal bath:

\[ \frac{dn_{\tilde{G}}}{dt} + 3Hn_{\tilde{G}} = C(T)n_{\text{rad}}^2. \tag{30} \]

Here \( H(T) \) is the Hubble parameter and \( n_{\text{rad}} = \frac{\zeta(3)}{\pi^2} T^3 \) is the number density of a relativistic bosonic degree of freedom. From eqs. (29) and (30) one obtains for the density of light gravitinos and the corresponding contribution to \( \Omega h^2 \) at temperatures \( T < 1 \text{ MeV} \), i.e. after nucleosynthesis,

\[ Y_{\tilde{G}} \simeq 3.2 \cdot 10^{-10} \left( \frac{T_B}{10^{10} \text{ GeV}} \right) \left( \frac{100 \text{ GeV}}{m_{\tilde{G}}} \right)^2 \left( \frac{m_{\tilde{g}}(\mu)}{1 \text{ TeV}} \right)^2, \tag{31} \]

\[ \Omega_{\tilde{G}} h^2 = m_{\tilde{G}} Y_{\tilde{G}}(T)n_{\text{rad}}(T)\rho_c^{-1} \simeq 0.60 \left( \frac{T_B}{10^{10} \text{ GeV}} \right) \left( \frac{100 \text{ GeV}}{m_{\tilde{G}}} \right)^2 \left( \frac{m_{\tilde{g}}(\mu)}{1 \text{ TeV}} \right)^2. \tag{32} \]

Here we have used \( g(T_B) = 0.85; \ \rho_c = 3H_0^2M^2 \) is the critical energy density, and \( m_{\tilde{g}}(T) = \frac{\theta(T)}{g^2(\mu)} m_{\tilde{g}}(\mu) \gg m_{\tilde{G}}, \) with \( \mu \sim 100 \text{ GeV}. \)

The primordial synthesis of light elements (BBN) yields stringent constraints on the amount of energy which may be released after nucleosynthesis by the decay of heavy nonrelativistic particles into electromagnetically and strongly interacting relativistic particles. These constraints have been studied in detail by several groups. Depending on the lifetime of the decaying particle \( X \) its energy density cannot exceed an upper bound. Sufficient conditions are

\begin{align*}
(I) \quad m_X Y_X(T) &< 4 \cdot 10^{-10} \text{ GeV}, \quad \tau < 2 \cdot 10^6 \text{ sec}, \tag{33} \\
(II) \quad m_X Y_X(T) &< 4 \cdot 10^{-12} \text{ GeV}, \quad \tau \text{ arbitrary}, \tag{34}
\end{align*}

where \( Y_X(T) = n_X(T)/n_{\text{rad}}(T). \)

Gravitinos interact only gravitationally. Hence, their existence leads almost unavoidably to a density of heavy particles which decay after nucleosynthesis. The partial width for the decay of an unstable gravitino into a gauge boson \( B \) and a bino \( \tilde{b} \) is given by \( (m_{\tilde{b}} \ll m_{\tilde{G}}) \),

\[ \Gamma(\tilde{G} \to B\tilde{b}) \simeq \frac{1}{32\pi} \frac{m_{\tilde{G}}^3}{M^2} \simeq 4 \cdot 10^8 \left( \frac{100 \text{ GeV}}{m_{\tilde{G}}} \right)^3 \text{ sec}^{-1}. \tag{35} \]
Figure 6: Neutralino relic density versus neutralino mass. The horizontal lines bound the region $0.025 < \Omega_\chi h^2 < 1$.

If for a fermion $\psi$ the decay into a final state with a scalar $\phi$ in the same chiral multiplet and a gravitino is kinematically allowed, the partial width reads ($m_\psi \gg m_\phi$),

$$\Gamma(\psi \rightarrow \tilde{G}\phi) = \Gamma(\psi \rightarrow \tilde{G}\phi^*) \simeq \frac{1}{96\pi} \frac{m_\psi^5}{m_\tilde{G}^2 M^2}.$$ (36)

Given these lifetimes and the mass spectrum of superparticles in the MSSM one can examine whether one of the conditions (I) and (II) on the energy density after nucleosynthesis is satisfied.

Consider first a typical example of supersymmetry breaking masses in the MSSM, $m_b < m_\tilde{G} \simeq 100 \text{ GeV} < m_3 \simeq 500 \text{ GeV}$, and $T_B \simeq 10^{10} \text{ GeV}$. From eqs. (31) and (35) we conclude $\tau_\tilde{G} \simeq 4 \cdot 10^8 \text{ sec}$, $m_\tilde{G} Y_\tilde{G}(T) \simeq 4 \cdot 10^{-9} \text{ GeV}$. According to condition (II) (34) this energy density exceeds the allowed maximal energy density by 3 orders of magnitude. This clearly illustrates the ‘gravitino problem’! The stringent constraints from BBN can be evaded for very light gravitinos, with $m_\tilde{G} < 1 \text{ keV}$ and possibly also for very heavy gravitinos, with $m_\tilde{G} = \mathcal{O}(10 \text{ TeV})$.

Another interesting possibility is the case where the gravitino is the LSP with a mass $\mathcal{O}(100 \text{ GeV})$. In this case one has to worry about the decays of the next-to-lightest superparticle (NSP) after nucleosynthesis. The lifetime constraint of condi-
Figure 7: Neutralino composition $Z_g/(1-Z_g)$ versus neutralino mass for $0.025 < \Omega_\chi h^2 < 1$.

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Figure 8: Upper and lower bounds on the NSP mass as function of the gravitino mass. The full lines represent the upper bound on the gluino mass $m_{\tilde{g}}(\mu) > m_{\text{NSP}}$ for different reheating temperatures. The dashed line is the lower bound on $m_{\text{NSP}}$ which follows from the NSP lifetime. A higgsino-like NSP with a mass in the shaded area satisfies all cosmological constraints including those from primordial nucleosynthesis.

this mass range and with a lifetime $\tau < 2 \cdot 10^6$ sec are compatible with the constraints from primordial nucleosynthesis. Note that this is a sufficient, yet not necessary condition for satisfying the bound $\Omega h^2 < 0.008$. Very small neutralino densities are also obtained for other sets of MSSM parameters.

Finally, we have to discuss the necessary condition that gravitinos do not overclose the universe,

$$\Omega_{\tilde{G}} h^2 < 1.$$  

(38)

Since $m_{\tilde{G}} < m_{\chi} < m_{\tilde{g}}$, the gravitino density is given by eq. (32). Hence, the condition (38) yields an upper bound on the NSP mass $m_{\chi}$ which depends on the gravitino mass and the baryogenesis temperature. The different constraints are summarized in fig. 8, which illustrates that for a wide range of MSSM parameters, where

$$m_{\tilde{G}} < m_{\chi} < m_{\tilde{g}},$$  

(39)

and $80 \text{ GeV} < m_{\chi} < 300 \text{ GeV}$, the baryogenesis temperature may be as large as $O(10^{10})$ GeV.
Figure 9: Contribution of gravitinos to the density parameter $\Omega h^2$ for different gravitino masses $m_{\tilde{G}}$ as function of the reheating temperature $T_B$. The gluino mass has been set to $m_{\tilde{g}}(\mu) = 500$ GeV.

It is remarkable that for temperatures $T_B = 10^8 \ldots 10^{11}$ GeV, which are natural for leptogenesis, and for gravitino masses in the range $m_{\tilde{G}} = 10^0 \ldots 10^3$ GeV, which is expected for gravity induced supersymmetry breaking, the relic density of gravitinos is cosmologically important (cf. fig. 9). As an example, for $T_B \simeq 10^{10}$ GeV, $m_{\tilde{g}}(\mu) \simeq 500$ GeV, and $m_{\tilde{G}} \simeq 50$ GeV, one has $\Omega_{\tilde{G}} h^2 \simeq 0.30$.

4. Summary

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 $\Lambda_{\text{GUT}} \sim 10^{16}$ GeV, the observed baryon asymmetry $n_B/s \sim 10^{-10}$ is naturally explained by the decay of heavy Majorana neutrinos. The corresponding baryogenesis temperature is $T_B \sim 10^{10}$ GeV.

In supersymmetric models implications for the mass spectrum of superparticles can be derived. The rather large baryogenesis temperature leads to a high density of gravitinos. Depending on the masses of the other superparticles their late decay may change the primordial abundances of light elements in disagreement with observation. This ‘gravitino problem’ can be avoided for very light gravitinos, with $m_{\tilde{G}} < 1$ keV, and possibly also for very heavy gravitinos, with $m_{\tilde{G}} = \mathcal{O}(10 \text{ TeV})$. Another interesting possibility is that the gravitino is the LSP, with a higgsino as NSP. It is intriguing that for a mass range $m_{\tilde{G}} = 10^{0} \ldots 10^{2}$ GeV and reheating temperatures $T_B = 10^{8} \ldots 10^{11}$ GeV, which naturally occur after inflation, one obtains a gravitino contribution to cold dark matter with $\Omega h^2 = \mathcal{O}(1)$.

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