Leptogenesis

Michael Plümacher

Theory Division, CERN, CH-1211 Geneva 23, Switzerland

Abstract. Thermal leptogenesis explains the baryon asymmetry of the universe by the out-of-equilibrium decays of heavy right-handed neutrinos. In the minimal seesaw model this leads to interesting implications for light neutrino properties. In particular, quasi-degenerate light neutrino masses are incompatible with leptogenesis. An upper bound on light neutrino masses of 0.2 eV can be derived, which will be tested by forthcoming laboratory experiments and cosmology.

1. Introduction

The observed baryon asymmetry of the universe is one of the most intriguing problems of particle physics and cosmology. The precision of measurements of the asymmetry has significantly improved with the observation of acoustic peaks in the cosmic microwave background (CMB). The BOOMERanG and DASI experiments have measured the baryon asymmetry with a (1σ) standard error of $\sim 15\%$ [1, 2],

$$\eta_{B0}^{CMB} = (6.0^{+1.1}_{-0.8}) \times 10^{-10} \;, \tag{1}$$

which is consistent with the results of standard primordial nucleosynthesis (BBN), obtained from the measurements of relic nuclear abundances.

A simple and elegant explanation of the observed baryon asymmetry is offered by neutrino physics. During the past years data on atmospheric and solar neutrinos have provided strong evidence for neutrino masses and mixings. In the seesaw mechanism [3] the smallness of these neutrino masses is naturally explained by the mixing of the left-handed neutrinos with heavy Majorana neutrinos. Further, the connection between baryon and lepton number in the high-temperature, symmetric phase of the standard model due to rapid sphaleron transitions [4] is by now firmly established [5]. As in classical GUT baryogenesis [6], out-of-equilibrium decays of the heavy Majorana neutrinos can then generate a lepton asymmetry which, by sphaleron processes, is partially transformed into a baryon asymmetry [7].

A beautiful aspect of this 'leptogenesis' mechanism is the connection between the cosmological baryon asymmetry and neutrino properties. This connection is established by standard kinetic calculations [8, 9, 10], very much like in big bang nucleosynthesis [6], where light nuclei play the role analogous to leptons in leptogenesis. The requirement of 'successful baryogenesis', i.e. the existence of neutrino masses and mixings for which the predicted and the observed value of the baryon asymmetry are in agreement, constitutes a severe test for models of neutrino masses, which has been extensively explored during the past years [11].

The dynamical generation of a baryon asymmetry requires that the particle interactions do not conserve baryon number, C and CP. In leptogenesis these conditions are realized by the couplings of the heavy Majorana neutrinos N_i. Their decays can generate an asymmetry in the number of leptons and antileptons, and therefore in $B-L$. The crucial departure from thermal equilibrium is provided by the expansion of the universe. At temperatures $T \sim \mathcal{O}(M_i)$ the abundance of heavy neutrinos exceeds the thermal abundance due to their weak interactions with the thermal bath.

In the following we shall assume that the dominant contribution to the baryon asymmetry is given by decays of N_1, the lightest of the heavy Majorana neutrinos. This assumption is well justified in the case of a mass hierarchy among the heavy neutrinos, i.e. $M_1 \ll M_2, M_3$, and it is also known to be a good approximation, if $M_{2,3} - M_1 = \mathcal{O}(M_1)$.

A quantitative description of the generation of a baryon asymmetry is obtained by means of two kinetic equations, one for the number of heavy neutrinos N_1 and one for the $B-L$ asymmetry [8, 9, 10]:

$$\frac{dN_{N_1}}{dz} = -\,(D+S)\,(N_{N_1} - N_{N_1}^{\rm eq})\,, \tag{2}$$

$$\frac{dN_{B-L}}{dz} = -\,\varepsilon_1\,D\,(N_{N_1} - N_{N_1}^{\rm eq}) - W\,N_{B-L}\,, \tag{3}$$

where $z \equiv M_1/T$ and ε_1 is the CP asymmetry in the decay of the right-handed neutrinos. N_i is the particle number in a comoving volume element $R_\star^3(t)$, chosen to contain one photon at a time $t_\star$ before the onset of leptogenesis. For a massless boson with g_B degrees of freedom in thermal equilibrium one has $N_B = g_B/2$, whereas for a fermion with g_F degrees of freedom $N_F = 3g_F/8$. The final baryon asymmetry is conveniently expressed in terms of the baryon-to-photon ration $\eta_{B0} = N_B(z)/N_\gamma(z)$. In the simple case of isentropic expansion of the universe one obtains $\eta_{B0} \simeq 0.036 N_{B0}$.

The relevant processes in the thermal plasma contributing to the equations are: decays and inverse decays of N_1 into leptons and Higgs bosons $N_1 \leftrightarrow l\phi$, and anti-leptons and anti-Higgs bosons, $N_1 \leftrightarrow \bar{l}\bar{\phi}$, $\Delta L = 1$ scatterings of N_1 off top-quarks and $\Delta L = 2$ processes mediated by the exchange of all heavy right-handed neutrinos. The first three contribute all together to modify the N_1 abundance. Denoting thermally averaged reaction rates by Γ_i, the term $D \equiv \Gamma_D/(H\,z)$ accounts for the decays and inverse decays, where H is the Hubble parameter. Similarly, the term $S \equiv \Gamma_S/(H\,z)$ accounts for the $\Delta L = 1$ scatterings. The decays are also the source term for the generation of the $B-L$ asymmetry, the first term in the second equation, while all the other processes contribute to the washout term $W \equiv \Gamma_W/(H\,z)$ that competes with the decay source term.

The contribution to W from $\Delta L = 2$ scatterings can conveniently be separated into two parts. The first part, $W_{\rm res}$, comes from the resonance contribution from N_1, which typically dominates in the temperature range from $z \simeq 1$ down to $z \simeq 10$. The remaining part, $W_{\mathrm{non-res}}$, is dominant at low temperatures, $z \lesssim 10$, where the scattering amplitudes with N_i exchange are proportional to the light neutrino mass matrix m_ν.

In order to understand the dependence of the solutions of the rate equations (2) and (3) on the neutrino parameters, it is crucial to note that the reaction rates are all dimensionless functions of the inverse temperature z, proportional to

$$D,\ S,\ W_{\mathrm{res}} \propto \frac{M_{\mathrm{Pl}}\widetilde{m}_1}{v^2}\,, \quad W_{\mathrm{non-res}} \propto \frac{M_{\mathrm{Pl}}M_1\,\overline{m}^2}{v^4}\,, \tag{4}$$

where v is the vacuum expectation value of the standard model Higgs field. The effective neutrino mass $\widetilde{m}_1$ [9] is given by

$$\widetilde{m}_1 = \frac{(m_D^\dagger m_D)_{11}}{M_1}\,, \tag{5}$$

where m_D is the Dirac neutrino mass matrix. Further, $\overline{m}^2 = \mathrm{tr}(m_\nu^\dagger m_\nu)$, is the quadratic mean of light neutrino masses, up to a factor $1/\sqrt{3}$. Eq. (4) implies that, as long as $W_{\mathrm{non-res}}$ can be neglected, the generated lepton asymmetry is independent of M_1.

2. Constraints on neutrino parameters

In this section we shall compare the predicted baryon asymmetry with the value measured from CMB (cf. eq. (1)). From eq. (3) it is easy to see that the final asymmetry N_{B-L} has to depend linearly on the CP asymmetry ε_1:

$$N_{B-L} = -\frac{3}{4}\varepsilon_1\kappa_0\,. \tag{6}$$

Here we have neglected a possible initial $B-L$ asymmetry present before leptogenesis. This usually is a good aproximation, since $\Delta L = 2$ processes mediated by heavy neutrinos are strong enough at high temperatures, $z \ll 1$, to washout an initial asymmetry.

The efficiency factor κ_0 is normalized to be 1 in the limiting case where an initial thermal abundance of heavy neutrinos decays far out of equilibrium, where all washout scattering processes can be neglected, i.e. in the case $W = 0$. In thermal leptogenesis this yields the maximal asymmetry, and one always has $\kappa_0 \leq 1$. In general, κ_0 depends only on three parameters: the effective neutrino mass $\widetilde{m}_1$, the quadratic mean of light neutrino masses $\overline{m}$ and the heavy neutrino mass M_1.

In Fig. 1 we have plotted the predicted present baryon asymmetry as a function of the parameter $\widetilde{m}_1$ for different values of M_1, assuming a typical value of the *CP* asymmetry, $\varepsilon_1 = -10^{-6}$ and $\overline{m} = 0.05\,\mathrm{eV} \simeq \sqrt{\Delta m^2_{\mathrm{atm}}}$, which corresponds to the case of hierarchical neutrinos. We have performed the calculations both for an initially vanishing N_1 abundance (thick lines) and for a thermal initial N_1 abundance, $N_{N_1}^{\mathrm{in}} = 3/4$ (thin lines).

In Fig. 2 we plot the dependence of the baryon asymmetry on the parameters $\widetilde{m}_1$ and M_1 by means of iso-κ_0 curves of the efficiency factor, again for the choice $\overline{m} = 0.05\,\mathrm{eV}$. For vanishing initial N_1 abundance the enclosed domains have a finite extension in $\widetilde{m}_1$; for thermal initial abundance there is no boundary at small $\widetilde{m}_1$.

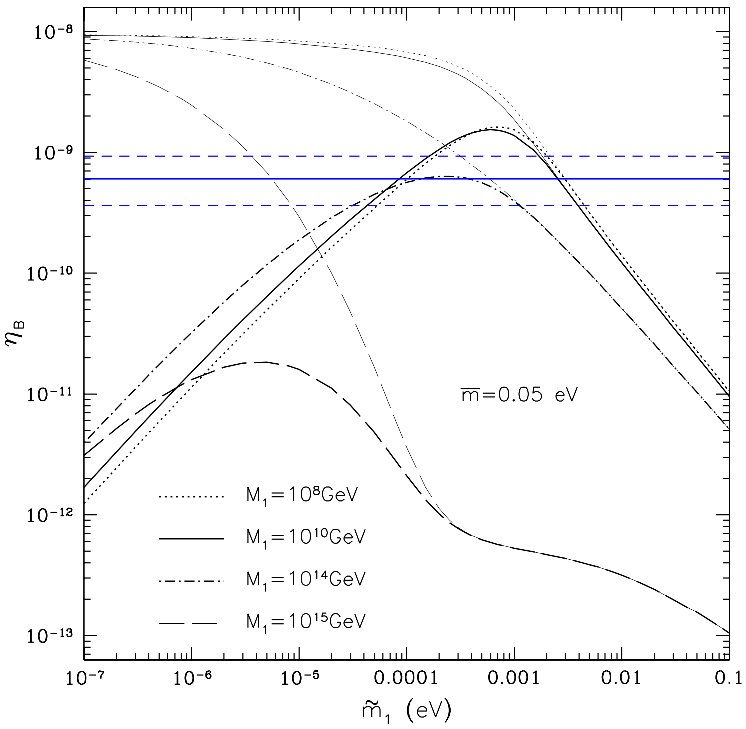

Figure 1. Predicted baryon-to-photon ratio η_{B0} as a function of $\widetilde{m}_1$, for $\varepsilon_1 = -10^{-6}$ and for the indicated values of M_1. The horizontal solid and dashed lines indicate the mean and the upper/lower values (3σ) of η_{B0}^{CMB}, respectively (see eq. (1)).

Figures 1 and 2 clearly show the existence of two different regimes: the domain of 'small' N_1 masses, $M_1 < 10^{13}\,\mathrm{GeV}(0.1\,\mathrm{eV}/\overline{m})^2$, and the domain of 'large' N_1 masses, $M_1 > 10^{13}\,\mathrm{GeV}(0.1\,\mathrm{eV}/\overline{m})^2$. The dependence of the boundary on $\overline{m}$ is determined by the behaviour of the non-resonant washout rate, $W_{\mathrm{non-res}} \propto M_1\overline{m}^2$.

In the small M_1 regime the non-resonant part $W_{\mathrm{non-res}}$ of the washout rate is negligible to a very good approximation, and there is only a small dependence on M_1 and $\overline{m}$ in the final asymmetry. Hence, the efficiency factor κ_0 depends approximately only on $\widetilde{m}_1$. For an initially vanishing N_1 abundance the suppression of the final asymmetry for small $\widetilde{m}_1$ is due to the fact that Yukawa interactions are not strong enough to create an initial thermal distribution of right-handed neutrinos. On the other hand, assuming an initial thermal abundance, the neutrinos decay far out of equilibrium and one approaches the limit $\kappa_0 \to 1$.

For large $\widetilde{m}_1$, the Yukawa interactions become strong enough to create an initial thermal distribution of heavy neutrinos, i.e. one becomes independent of initial conditions. The suppresion of the baryon asymmetry is now due to the washout processes W.

In the large M_1 regime the non-resonant washout rate $W_{\mathrm{non-res}}$ dominates, leading to an exponential suppression of the produced baryon asymmetry. Here the iso-κ_0 curves become almost independent of $\widetilde{m}_1$. Hence, in the most general case, the final baryon

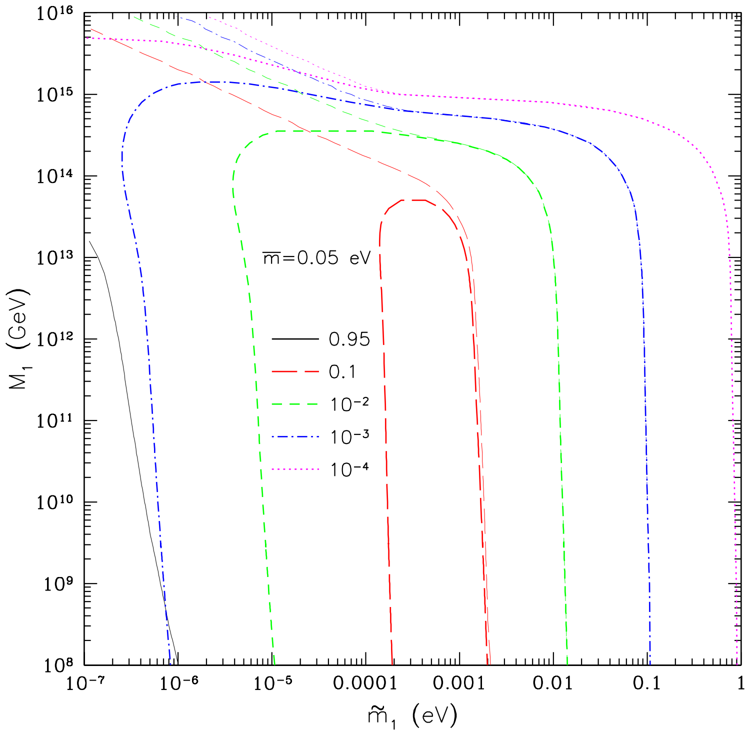

Figure 2. Iso-κ_0 curves for $N_{N_1}^{\rm in} = 0$ (thick lines) and $N_{N_1}^{\rm in} = 3/4$ (thin lines) assuming $\overline{m} = 0.05$ eV $\simeq \sqrt{\Delta m^2_{\rm atm}}$.

asymmetry depends on only four parameters: ε_1, $\widetilde{m}_1$, M_1 and $\overline{m}$.

3. The maximal baryon asymmetry

Up to now we have treated the CP asymmetry ε_1 as an independent parameter. However, it can be shown that the CP asymmetry satisfies an upper bound [12]

$$|\varepsilon_1| < \varepsilon(M_1, \overline{m}) = \frac{3}{16\pi}\frac{M_1}{v^2}\frac{\Delta m^2_{\rm atm} + \Delta m^2_{\rm sol}}{m_3} \, . \tag{7}$$

The maximal baryon asymmetry produced by N_1 decays is therefore given by:

$$\eta_{B0}^{\rm max}(\widetilde{m}_1, M_1, \overline{m}) \simeq 10^{-2}\,\varepsilon(M_1, \overline{m})\,\kappa_0(\widetilde{m}_1, M_1, \overline{m}) \, . \tag{8}$$

$\eta_{B0}^{\rm max}$ has to be larger than the observed baryon asymmetry (cf. eq. (1)) for this mechanism to be succesful. This yields a lower bound on M_1 [12]. For zero (thermal) initial N_1 abundance one obtains [10],

$$M_1 \gtrsim 2.4\,(0.4) \times 10^9\,{\rm GeV}\left(\frac{2.5\times 10^{-3}{\rm eV}^2}{\Delta m^2_{atm}}\right)^{1/2} \, . \tag{9}$$

This is illustrated in Fig. 3, where iso-$\eta_{B0}^{\rm max}$ curves in the $(\widetilde{m}_1, M_1)$-plane are shown for $\overline{m} = 0.05$ eV. For values $\widetilde{m}_1$ and M_1 enclosed by these curves the baryon asymmetry η_{B0}^{CMB} can be obtained. In the case of $N_{N_1}^{\rm in} = 0$, the allowed regions are closed domains; for $N_{N_1}^{\rm in} = 3/4$ there is no lower limit on $\widetilde{m}_1$.

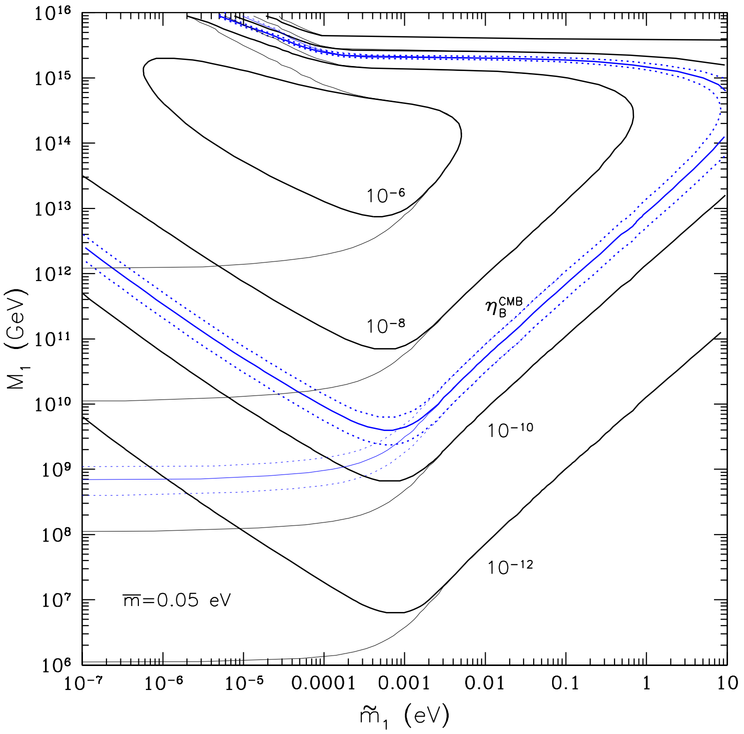

Figure 3. Iso-$\eta_{B0}^{\rm max}$ curves for $N_{N_1}^{\rm in} = 0$ (thick lines) and $N_{N_1}^{\rm in} = 3/4$ (thin lines) for the case of hierarchical light neutrino masses ($\overline{m} = 0.05$ eV $\simeq \sqrt{\Delta m^2_{\rm atm}}$). The region within the dashed lines is currently allowed by the CMB constraint (cf. eq. (8)) for the (3σ) lower η_{B0}^{CMB} value.

4. A bound on light neutrino masses

It can be shown that the effective neutrino mass $\widetilde{m}_1$ is bounded from below by the lightest neutrino mass m_1 and from above by m_3 [13, 10]. Hence, for large $\overline{m}$, corresponding to quasi-degenerate light neutrinos, one has $\widetilde{m}_1 \simeq \overline{m}$.

Increasing $\overline{m}$ leads to a suppressed baryon asymmetry, since the bound (7) on the CP asymmetry becomes stronger, and because the washout rates become more effective. Hence, there is a maximal value of $\overline{m}$ above which the observed baryon asymmetry cannot be produced in N_1 decays. This is illustrated in Fig. 4, where we have displayed the curves defined by $\eta_{B0}^{\rm max}(\widetilde{m}_1, M_1, \overline{m}) = \eta_{B0}^{CMB}$ and by $\widetilde{m}_1 = m_1(\overline{m})$, for $\overline{m} = 0.25$ eV (dashed), 0.29 eV (dotted) and 0.30 eV (full). For the two smaller values of $\overline{m}$ domains exist which allow large enough baryon asymmetry. For $\overline{m} = 0.30$ eV successful baryogenesis is no longer possible [14]. In terms of light neutrino masses this corresponds to the limits $m_3 < 0.18\,$eV and $m_1 \simeq m_2 < 0.17\,$eV.

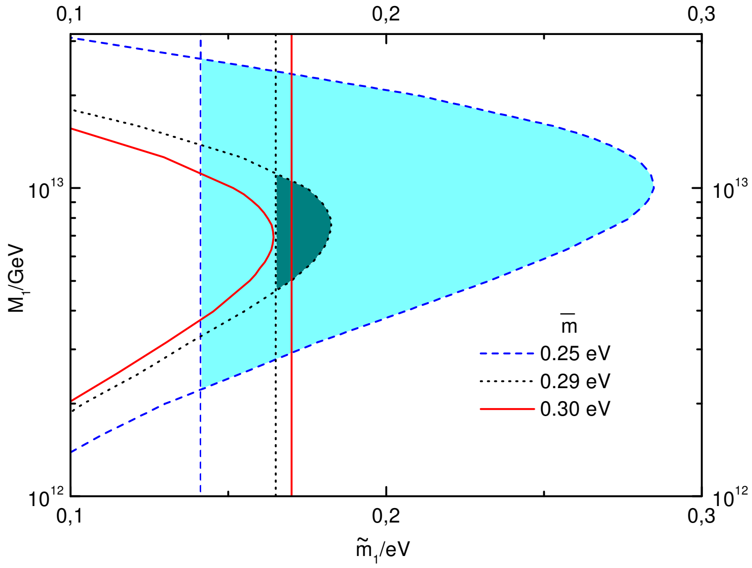

Figure 4. Regions in the $\widetilde{m}_1$-M_1-plane allowed by baryogenesis for different values of $\overline{m}$. The boundaries are defined by $\eta_{B0}^{\max}(\widetilde{m}_1, M_1, \overline{m}) = \eta_{B0}^{CMB}$ and by $\widetilde{m}_1 = m_1(\overline{m})$.

5. Conclusions

We have shown that in the minimal version of thermal leptogenesis, where the lightest of the heavy Majorana neutrinos N_1 is the source of the baryon asymmetry, the final baryon asymmetry is determined by only four parameters: the CP asymmetry ε_1, the heavy neutrino mass M_1, the effective light neutrino mass $\widetilde{m}_1$ and the quadratic mean $\overline{m}$ of light neutrino masses.

The observed baryon asymmetry can only be explained if M_1, and therefore the baryogenesis temperature $T_B \sim M_1$, is above 2×10^9 GeV and if light neutrino masses are not too large, $m_3 < 0.18\,\mathrm{eV}$ and $m_1 \simeq m_2 < 0.17\,\mathrm{eV}$. Using the unitarity of the leptonic mixing matrix, $\sum_i |U_{ei}|^2 = 1$, one then obtains for the mass measured in tritum β-decay, $m_{\nu_e} < 0.18$ eV. This value is smaller than the sensitivity aimed at by the KATRIN collaboration [15]. The same bound as for m_{ν_e} one also finds for the Majorana mass measured in neutrinoless double β-decay, $m_{ee} < 0.18$ eV. This prediction can be tested by the GENIUS project [16]. Cosmological bounds on neutrino masses will also improve, and the Sloan Digital Sky Survey combined with CMB data from the MAP experiment is expected to reach the sensitivity [17] of the leptogenesis bound, $\sum_i m_i \leq 0.52$ eV.

Acknowledgments

I would like to thank W. Buchmüller and in P. Di Bari for an enjoyable collaboration on the topic of this review. Further, I would like to thank the organizers of NuFact'02 fo organizing a stimulating and fruitful meeting.

[1] P. de Bernardis et al., Astrophys. J. **564** (2002) 559

[2] C. Pryke et al., Astrophys. J. **568** (2002) 46

[3] 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

[4] V. A. Kuzmin, V. A. Rubakov, M. E. Shaposhnikov, Phys. Lett. **B 155** (1985) 36

[5] D. Bödeker, G. D. Moore, K. Rummukainen, Phys. Rev. **D 61** (2000) 056003

[6] E. W. Kolb, M. S. Turner, *The Early Universe*, Addison-Wesley, New York, 1990

[7] M. Fukugita, T. Yanagida, Phys. Lett. **B 174** (1986) 45

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

[9] M. Plümacher, Z. Phys. **C 74** (1997) 549

[10] W. Buchmüller, P. Di Bari and M. Plümacher, Nucl. Phys. B **643** (2002) 367

[11] For a review and references, see
W. Buchmüller, M. Plümacher, Int. J. Mod. Phys. **A 15** (2000) 5047

[12] S. Davidson and A. Ibarra, Phys. Lett. B **535** (2002) 25

[13] M. Fujii, K. Hamaguchi and T. Yanagida, Phys. Rev. D **65** (2002) 115012

[14] W. Buchmüller, P. Di Bari and M. Plümacher, Phys. Lett. B **547** (2002) 128

[15] KATRIN Collaboration, A. Osipowicz et al., hep-ex/0109033

[16] H. V. Klapdor-Kleingrothaus, Nucl. Phys. Proc. Suppl. **100** (2001) 350

[17] W. Hu, D. J. Eisenstein, M. Tegmark, Phys. Rev. Lett. **80** (1998) 5255