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Baryogenesis in the νMSM: recent developments*

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Abstract. Extension of the Standard Model with two right-handed neutrinos provides an economic and testable explanation of the origin of the baryon asymmetry of the Universe. We review recent progress in understanding dynamics of the asymmetry generation. We also present results of a new study of the parameter space of the model. These results demonstrate that the region of the parameter space in which the observed value of baryon asymmetry can be reproduced is larger than it was previously obtained.

1 Introduction

The neutrino oscillations are among the few well established phenomena beyond the Standard Model (SM). Another one is the baryon asymmetry of the Universe (BAU). Both of them can be explained within the framework of the νMSM, that is the extension of the SM with three right-handed (RH) Majorana neutrinos with masses below the electroweak scale. The left-handed neutrinos obtain their masses via the see-saw mechanism [3–8], whereas the BAU is generated in CP-violating oscillations of the RH neutrinos [9, 10]. Potential testability of the model and it's minimality have attracted a lot of attention from both theoretical and experimental sides. An incomplete list of theoretical works includes [1, 11–33]. New direct experimental searches for the RH neutrinos were carried out as well [34–39].[1] There are several new proposals of the experiments which will be very sensitive to these RH neutrinos, such as NA62 [41], SHiP [42] and MATHUSLA [43].

In this talk we discuss a procedure of an accurate calculation of the BAU in the νMSM and other models of low-scale leptogenesis [1]. Then we list other recent developments in the subject. Finally, we present the results of the new study of the parameter space of the model [2].

2 Baryogenesis with GeV-scale RH neutrinos

In this section we briefly describe the model with two RH neutrinos, which can be viewed as a part of the νMSM. Two RH neutrinos, N_2 and N_3, are enough to reproduce the oscillation data by means of the see-saw mechanism since there are two measured mass differences. The

*The talk is based on refs. [1, 2].

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[1]It was shown that indirect searches are incapable of finding RH neutrinos of the νMSM [40].

first RH neutrino N_1 can play a role of the Dark Matter particle [44–48]. However, Lyman-α and X-ray constraints put strong bounds on it's Yukawa couplings [49]. Therefore, N_1 is effectively decoupled from the model. The Lagrangian of the νMSM reads

$$\mathcal{L} = \mathcal{L}_{SM} + \bar{N}_I \gamma^\mu \partial_\mu N_I - F_{\alpha I} \bar{L}_\alpha \tilde{\Phi} N_I - \frac{M_{IJ}}{2} \bar{N}_I^c N_J + h.c. \tag{1}$$

where $\mathcal{L}_{SM}$ is the SM part, $F_{\alpha I}$ is the matrix of Yukawa couplings, L_α are the left-handed lepton doublets and Φ is the Higgs doublet. We assume that the charged lepton Yukawa couplings and the Majorana mass term for the RH neutrinos M_{IJ} are diagonal.

In order to reproduce the experimentally observed values of active neutrino mass differences and mixings Casas and Ibarra suggested to decompose the matrix of Yukawa couplings in the following way [50].

$$F = \frac{i}{v_0} U^{PMNS} m_\nu^{1/2} \Omega m_N^{1/2}, \tag{2}$$

where U^{PMNS} is the PMNS matrix and the complex orthogonal matrix Ω controls the misalignment of the Yukawas. In the case of the RH neutrinos it can be chosen in the following form.

$$\Omega = \begin{pmatrix} 0 & 0 \\ \cos\omega & \sin\omega \\ -\xi \sin\omega & \xi \cos\omega \end{pmatrix} \quad \text{for NH,} \tag{3}$$

$$\Omega = \begin{pmatrix} \cos\omega & \sin\omega \\ -\xi \sin\omega & \xi \cos\omega \\ 0 & 0 \end{pmatrix} \quad \text{for IH,} \tag{4}$$

where ω is a complex angle. Throughout this text the abbreviations NH and IH refer to the normal and inverse mass hierarchy respectively.

In the two RH case, the Majorana mass matrix in (1) can be parametrized by a common mass M and a mass splitting ΔM. If only two active neutrinos are massive, the PMNS matrix contains two CP phases, a Dirac one, δ, and a Majorana phase, denoted by η. At the end we have six free parameters of the model. They are listed, along with considered ranges, in table 1. Let us stress that for all values of the parameters listed in table 1 one reproduces the oscillation data.

M, GeV	$\log_{10}(\Delta M/\text{GeV})$	$\text{Im}\,\omega$	$\text{Re}\,\omega$	δ	η
$[0.1 - 10]$	$[-17, -1]$	$[-7, 7]$	$[0, 2\pi]$	$[0, 2\pi]$	$[0, 2\pi]$

Table 1. Parameters of the theory and the ranges of these parameters which were considered in this work.

3 Freeze-out of the baryon number in low scale leptogenesis models

In the framework of leptogenesis the baryon asymmetry is reprocessed from the lepton asymmetry by the electroweak sphalerons [51, 52]. The total baryon and lepton numbers (B and $L = \sum L_\alpha$, $\alpha = e, \mu, \tau$) are related by the

$$B(T) = \chi(T) \sum_\alpha \Delta_\alpha(T), \tag{5}$$

where we defined $\Delta_\alpha = L_\alpha - B/3$ (this combination is not affected by sphalerons) and $\chi(T)$ is a coefficient of order of 1; for instance, in the symmetric phase it is equal to $-28/79$ [53].

The equilibrium formula (5) is valid as long as the sphaleron rate exceeds the rate of the lepton asymmetry generation. This is the case for high-scale leptogenesis, $T \gg 10^8$ GeV [54]. However, this assumption is violated in the low-scale leptogenesis. At temperatures around $T_{sph} \simeq 131.7$ GeV (the most recent calculation has been done in ref [55]) sphalerons become inoperative, but the lepton asymmetry is still being generated. Therefore, the actual amount of the BAU reproduced from the lepton asymmetry by sphalerons can be different from the value predicted by eq. (5) at $T = T_{Sph}$. In order to study this effect we will need to solve a system of kinetic equations describing the asymmetry generation plus an extra equations for the baryon number.

The production of lepton asymmetry is described by a set of kinetic equations. In the Higgs phase, kinetic equations generically can be written as:

$$\begin{aligned} \dot{n}_{\nu_\alpha} &= f_\alpha(\mathbf{n}_N, n_{\nu_\alpha}), \\ \dot{\mathbf{n}}_N &= g(\mathbf{n}_N, n_{\nu_\alpha}), \end{aligned} \tag{6}$$

where n_{ν_α} ($\alpha = e, \mu, \tau$) are the asymmetries of the number densities of left-handed neutrinos, $\mathbf{n}_N$ is a matrix of number densities and correlations of RH neutrinos.

This system is far from being realistic: (i) charged fermions are missing, (ii) sphaleron processes, which are fast at temperatures above $T_{sph} \simeq 131.7$ GeV [55], are not accounted for.

At temperatures of the low-scale leptogenesis all SM species are in equilibrium: $\mu_{\nu_\alpha} = \mu_{e_{L,\alpha}} = \mu_{e_{R,\alpha}} = \mu_\alpha$. So if one formulates the equations in terms of chemical potentials, the charged fermions will be automatically accounted for. Note that only n_{ν_α} are changing due to interactions with HNLs: $[\dot{n}_{\nu_\alpha}]_{HNLs} = [\dot{n}_\alpha]_{HNLs}$ Also $[\dot{n}_\alpha]_{HNLs} = [\dot{n}_{\Delta_\alpha}]_{HNLs}$, where $n_{\Delta_\alpha} = n_\alpha - n_B/3$. Since $B - L$ is preserved by sphalerons, we can write the l.h.s. of the equations in terms n_{Δ_α}.

$$\begin{aligned} \dot{n}_{\Delta_\alpha} &= f_\alpha(\mathbf{n}_N, \mu_\alpha), \\ \dot{\mathbf{n}}_N &= g_I(\mathbf{n}_N, \mu_\alpha). \end{aligned}$$

The neutrality of the electroweak plasma implies a non-trivial relation between the chemical potentials and the asymmetries [1, 56]

$$\mu_\alpha = \omega_{\alpha\beta}(T) n_{\Delta_\beta} + \omega_B(T) n_B, \tag{7}$$

where ω are the so-called susceptibility matrices. The matrices ω can be calculated by means of equilibrium thermodynamics [1]. Let us introduce the thermodynamical potential [57]

$$\begin{aligned} &\Omega(\mu, T) = \\ &\frac{1}{24}\Big(8T^2\mu_B^2 + 8T^2\mu_B\mu_Y + 6\mu_1^2T^2 + 6\mu_2^2T^2 + 6\mu_3^2T^2 + 22T^2\mu_T^2 + \\ &22T^2\mu_Y^2 - 8\mu_1T^2\mu_Y - 8\mu_2T^2\mu_Y - 8\mu_3T^2\mu_Y + 3\langle\Phi\rangle^2\mu_T^2 - 6\langle\Phi\rangle^2\mu_T\mu_Y + 3\phi^2\mu_Y^2\Big), \end{aligned}$$

where $\langle\Phi\rangle$ is the Higgs vacuum expectation value. The number densities of the conserved charges are

$$-\frac{\partial(\Omega/\mathcal{V})}{\partial\mu_\alpha} = n_\alpha, \quad -\frac{\partial(\Omega/\mathcal{V})}{\partial\mu_B} = n_B,$$

whereas the neutrality conditions read

$$\frac{\partial(\Omega/\mathcal{V})}{\partial\mu_Y} = 0, \quad \frac{\partial(\Omega/\mathcal{V})}{\partial\mu_T} = 0.$$

Solving equations above one calculates the susceptibility matrices entering eq. (7).

$$\omega(T) = \frac{1}{T^2}\begin{pmatrix} a & b & b \\ b & a & b \\ b & b & a \end{pmatrix}, \quad a = \frac{22\left(15x^2+44\right)}{9\,(17x^2+44)}, \quad b = \frac{8\left(3x^2+22\right)}{9\,(17x^2+44)}$$

and

$$\omega_B(T) = \frac{1}{T^2}\frac{4\left(27x^2+77\right)}{9\,(17x^2+44)},$$

where $x = \langle\Phi\rangle/T$.

In the same way one can derive a temperature dependent relation for the baryon number

$$n_{B^{eq}} = -\chi(T)\sum_\alpha n_{\Delta_\alpha}, \quad \chi(T) = \frac{4\left(27(\langle\Phi\rangle/T)^2+77\right)}{333(\langle\Phi\rangle/T)^2+869}. \tag{8}$$

The final ingredient is a separate kinetic equation for n_B [53, 58]

$$\dot{n}_B = -\Gamma_B(n_B - n_{B^{eq}}), \tag{9}$$

with

$$\Gamma_B = 9\frac{869+333(\langle\Phi\rangle/T)^2}{792+306(\langle\Phi\rangle/T)^2}\cdot\frac{\Gamma_{diff}(T)}{T^3}.$$

The Chern-Simons diffusion rate in a pure gauge theory [55]:

$$\Gamma_{diff} \simeq \begin{cases} T^4\cdot\exp\left(-147.7+0.83T/\mathrm{GeV}\right), & \text{broken phase,} \\ T^4\cdot 18\,\alpha_W^5, & \text{symmetric phase.}\end{cases}$$

We will solve the kinetic equations of ref. [31] along with the kinetic equation (9) to study the dynamics of decoupling of the baryon number. Let us introduce two approaches

1. **Approach 1**. A scenario of an instantaneous B freeze out.

   Baryon number density $n_B(T) = n_{B^{eq}}(T)$ for all temperatures above T_{sph} and $n_B(T) = n_{B^{eq}}(T_{sph})$ for all $T < T_{sph}$.

2. **Approach 2**. An approach with the separate kinetic equation for n_B.

   In this case one can follow the n_B during the freeze out, but at the cost of adding a new scale into the problem.

We found that the lepton asymmetry is the same in both approaches. We present two examples if figure 1. Still, the difference between BAU in the two approaches can be quite large for some values of the parameters. An example of a large deviation is shown in figure 2, $B^{\mathrm{appr.1}}/B^{\mathrm{appr.2}} \simeq -21.9$. So we can conclude that in some cases the assumption of the instantaneous freeze-out of sphalerons can lead to large errors. A numerical procedure which allows fast and accurate treatment of sphalerons has been suggested in ref [1].

4 A new study of the parameter space

As we have already mentioned, the νMSM has attracted a lot of attention in last years. The first study of the parameter space of the model has been performed in ref. [18]. However, a considerable progress in understanding of various rates entering the kinetic equations describing the asymmetry generation has been achieved since work since then. Let us list some of these developments

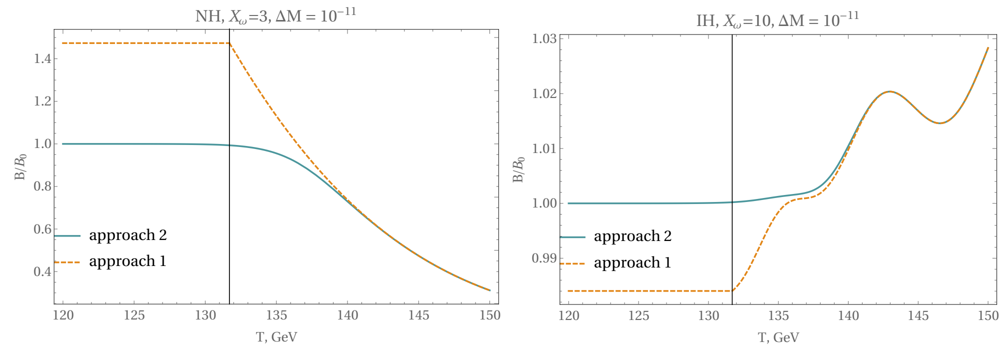

Figure 1. Y_B/Y_B^0 as a function of temperature in the approach with instant freeze-out (dotted orange line) and in the approach with the separate equation for B (blue line). Left panel: normal hierarchy, right panel: inverted hierarchy of neutrino masses. The thin black vertical line shows the sphaleron freeze-out temperature $T_{sph} = 131.7$ GeV. The common mass of HNLs is $M = 1$ GeV.

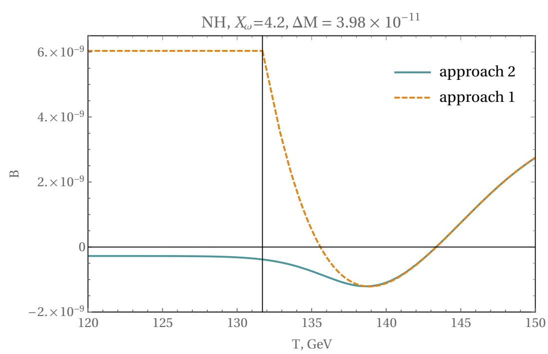

Figure 2. Y_B as a function of temperature. A case of large deviation. The black vertical line shows the sphaleron freeze out temperature $T_{sph} = 131.7$ GeV.

- Accurate computation of relevant rates [31, 56].
- Fermion number violating processes were included [30–32].
- Importance of the gradual freeze-out of sphalerons [1], see also the section above.

All this improvements have motivated us to perform a comprehensive study of the parameter space. We update the bounds on the parameters imposed by the requirement of successful baryogenesis [2].

The mixing of the RH neutrinos with active neutrinos is characterised by a function $|U|^2$, see, e.g. ref [18]. In figure 3 we present the new cosmological bounds on $|U|^2$. These bounds are obtained using the kinetic equations of refs [1, 31]. One can see, that the allowed region is significantly larger.

5 Conclusions

To summarize, we have discussed the importance of the gradual freeze-out of sphalerons in the models of low-scale leptogenesis. It has been shown that for some regims of the asymmetry generation the standard assumption of the instantaneous freeze-out of sphalerons

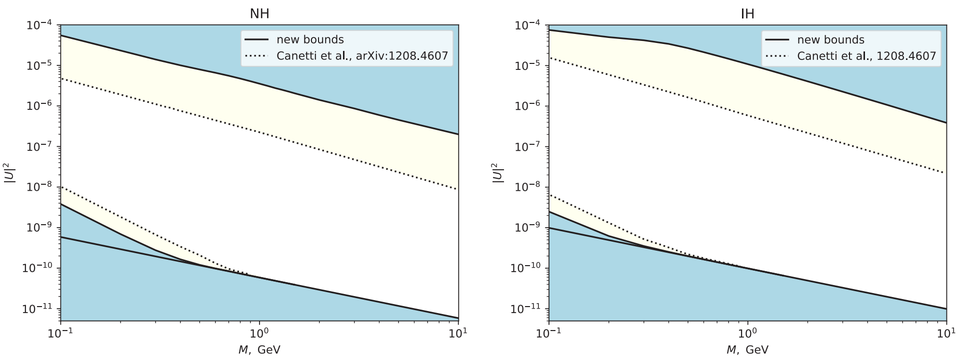

Figure 3. Bounds on the total mixing $|U|^2$ as functions of the common mass M. Within the white region the observed value of the BAU can be reproduced. The dotted line shows the bounds of the previous study [18].

can lead two large errors. Therefore, if one wants to ensure the accuracy for all values of the model parameters, it is more safe to solve a separate kinetic equation for a baryon number.

We have also demonstrated that the allowed region of the parameter space of the νMSM is larger than it was previously obtained. This opens new perspectives for direct detection experiments.

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