



# Dark matter and leptogenesis in gauged $B-L$ symmetric models embedding $\nu$ MSM

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## Abstract

We study the phenomenon of baryogenesis via leptogenesis in the gauged  $B-L$  symmetric models by embedding the currently proposed model  $\nu$ MSM. It is shown that the lightest right-handed neutrino of mass 100 GeV satisfy the leptogenesis constraint and at the same time representing a candidate for the cold dark matter. We discuss our results in parallel to the predictions of  $\nu$ MSM.

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## 1. Introduction

At present the atmospheric neutrino data [1] in the  $\nu_\mu-\nu_\tau$  oscillation and the solar neutrino data [2] in the  $\nu_e-\nu_\mu$  oscillation experiments are highly suggestive to believe small masses of the light neutrinos ( $\leq 1$  eV), either Dirac or Majorana. Assuming that the neutrinos are of Majorana type the small masses can be understood through the seesaw mechanism [3], which involves the right-handed neutrinos into the electroweak model, invariant under all the gauge transformations.

At the minimal cost we can add two right-handed neutrinos to the standard model (SM) Lagrangian to explain the tiny mass scales; the atmospheric neutrino mass ( $\Delta_{\text{atm}} = \sqrt{|m_3^2 - m_2^2|}$ ) and the solar neutrino mass ( $\Delta_{\text{sun}} = \sqrt{|m_2^2 - m_1^2|}$ ). However, in this scenario the seesaw mechanism gives rise to one of the light neutrino mass to be exactly zero. This is unwelcome if the neutrino masses are partially degenerate, albeit the hierarchical mass spectrum of the light neutrinos can be conspired in this scenario. Since the exact mass scales of the light neutrinos are

not known yet, we therefore add three right-handed neutrinos, gauge invariantly, to the SM Lagrangian.

In the thermal scenario the  $CP$  violating decay of the right-handed Majorana neutrinos can potentially explain the matter antimatter asymmetry [4], defined by

$$\frac{n_B}{n_\gamma} = 6.1 \times 10^{-10}, \quad (1)$$

of the present Universe as predicted by the Wilkinson Microwave Anisotropy Probe (WMAP) [5]. This requires the scale of operation of right-handed neutrinos to be  $\gtrsim 10^8$  GeV [6] and hence far beyond our hope to be verified in the near future accelerators.

An alternative is to consider the mechanisms which work at TeV scale [7–9]. In Ref. [7] it was proposed that the spontaneous breaking of the  $B-L$  gauge symmetry gives rise to a raw lepton asymmetry. The preservation of lepton asymmetry then requires a limited wash out through the lepton violating interactions mediated by the right-handed neutrinos and hence requiring the mass scale ( $M_1$ ) of lightest right-handed neutrino ( $N_1$ ) to be at the TeV scale or less. This needs to rethink whether these low mass scales of  $N_1$  can be compatible with the seesaw mechanism to give rise the Majorana mass matrix of the light

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neutrinos to be

$$m_\nu = -m_D^T M_R^{-1} m_D. \quad (2)$$

In particular, it was shown in Ref. [7] that a TeV mass of  $N_1$  is compatible with the seesaw if we assume that the Dirac mass matrix of the neutrinos is two orders less than that of charged leptons mass matrix.

Recently “ $\nu$ MSM” model has been proposed [10]. In this model the right-handed neutrinos are singlet under the SM gauge group. The mass of  $N_1$  in this case is constrained to

$$2 \leq M_1 \leq 5 \text{ keV}, \quad (3)$$

where the lower bound comes from the cosmic microwave background (CMB) and the matter power spectrum inferred from Lyman  $\alpha$  forest data [11] and the upper bound comes from the radiative decays of singlet right-handed neutrinos in dark matter (DM) halos limited by X-ray observations [12]. This requires, from Eqs. (2) and (3), that the Dirac Yukawa coupling  $h_\nu \sim 10^{-10}$  for  $m_\nu = 0.1 \text{ eV}$ . The tiny Yukawa coupling in this scenario makes the lightest right-handed neutrino decoupled from the thermal bath through out its evolution. However, this constraint is not applicable to  $N_2$  and  $N_3$ , the second and third generation of right-handed neutrinos. Hence they come into equilibrium through the large Yukawa couplings. This permits the authors in Ref. [13] to consider a mechanism to create a net lepton asymmetry in the right-handed neutrino sector through oscillations [13]. The net lepton asymmetry created in the right-handed sector is then transferred to the left-handed sector through the Yukawa coupling

$$\mathcal{L}_Y = (h_\nu)_{ij} \phi \bar{\psi}_{Li} N_j. \quad (4)$$

The lepton asymmetry is then transferred to baryon asymmetry through the nonperturbative sphaleron processes [14].

An important issue of the  $\nu$ MSM model is that the mass of  $N_1$  is severely constrained from the hot DM consideration. Further the lepton asymmetry produced by any mechanism other than the sterile neutrino oscillation will continue to survive and hence invalidating the leptogenesis constraints on the right-handed neutrino masses. Moreover the Dirac Yukawa couplings are very tiny.

As an alternative, in the present case we consider the low energy left–right symmetric model [15]. Since  $B-L$  is a gauge symmetry of the model any primordial asymmetry is erased. Further advantage of considering this model is that it can be easily embedded in the unified models like  $SO(10)$  or Pati–Salam and at the same time it can embed the “ $\nu$ MSM” by gauging the  $B-L$  symmetry. In this model, by assuming a normal mass hierarchy in the right-handed neutrino sector, we discuss the role of lightest right-handed neutrino in leptogenesis as well as for dark matter. It is shown that the lightest right-handed neutrino of mass 100 GeV can be a candidate for DM as well as satisfying the erasure constraint required for the preservation of lepton asymmetry.

Rest of the Letter is arranged as follows. In Section 2 we discuss the leptogenesis in the left–right symmetric models and then elucidate the possibility of bringing down the mass scale

of  $N_1$  to TeV scale or less. In Section 3 the constraint on the  $B-L$  breaking scale is discussed. In Section 4 we discuss the constraint on the mass scale of lightest right-handed neutrino from the flavor changing neutral current (FCNC). Section 5 is devoted to discuss the possibility of TeV scale right-handed neutrinos for the candidate of dark matter. Finally in Section 6 we summarize our results and put the conclusions.

## 2. Leptogenesis in gauged $B-L$ symmetric models and the possibility of TeV scale right-handed neutrino

In the following we consider the left–right symmetric model where  $B-L$  gauge symmetry emerges naturally. However, the arguments to be advocated below will remain valid as long as  $B-L$  is a gauge symmetry of the model.

### 2.1. Spontaneous CP-violation in $L-R$ symmetric model and leptogenesis

The main attraction of the left–right symmetric model lies in the lepton sector. The right-handed neutrinos, which were singlet under the SM gauge group  $SU(2)_L \times U(1)_Y$ , now non-trivially transforms under the left–right symmetric gauge group  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ . Since the right-handed neutrinos possess the  $B-L$  quantum number by one unit the Majorana mass can violate lepton number by two units and hence is a natural source of lepton asymmetry in the model.

The Higgs sector of the left–right symmetric model is very rich. It consists of two scalar triplets  $\Delta_L$  and  $\Delta_R$  which give Majorana masses to the right-handed neutrinos and a bidoublet  $\Phi$  which gives Dirac masses to the charged leptons and quarks. We assume that all the Yukawa couplings in the Higgs potential of the model are real. Thus the Lagrangian respects CP symmetry. The complex nature of the neutrino masses then come through the VEVs of the neutral Higgses [16]. In general there are four neutral Higgses in the model can potentially acquire VEVs and thereby breaking the left–right symmetry down to  $U(1)_{\text{em}}$ . Hence there are four phases associated with the neutral Higgses. However, the remnant global symmetry  $U(1)_L \times U(1)_R$  allows us to set two of the phases to zero. Therefore, only two of the phases have physical significance.

The breaking of left–right discrete symmetry in the early Universe gives rise to domain walls. It was shown in [17] that within the thickness of the domain walls the net CP violating phase becomes position dependent. Under these circumstances the preferred scattering of  $\nu_L$  over its CP-conjugate state ( $\nu_L^c$ ) produce a net raw lepton asymmetry [17]

$$\eta_L^{\text{raw}} \cong 0.01 v_w \frac{1}{g_*} \frac{M_1^4}{T^5 \Delta_w}, \quad (5)$$

where  $\eta_L^{\text{raw}}$  is the ratio of  $n_L$  to the entropy density  $s$ . In the right-hand side  $\Delta_w$  is the wall width and  $g_*$  is the effective thermodynamic degrees of freedom at the epoch with temperature  $T$ . Using  $M_1 = f_1 \Delta_T$ , with  $\Delta_T$  is the temperature dependent VEV acquired by the  $\Delta_R$  in the phase of interest, and

$\Delta_w^{-1} = \sqrt{\lambda_{\text{eff}}}\Delta_T$  in Eq. (5) we get

$$\eta_L^{\text{raw}} \cong 10^{-4} v_w \left(\frac{\Delta_T}{T}\right)^5 f_1^4 \sqrt{\lambda_{\text{eff}}}. \quad (6)$$

Here we have used  $g_* = 110$ . Therefore, depending on the various dimensionless couplings, the raw asymmetry may lie in the range  $O(10^{-4}-10^{-10})$ .

## 2.2. TeV scale right-handed neutrino and lepton asymmetry

In the previous section we saw that a net raw lepton asymmetry ( $\eta_L^{\text{raw}}$ ) is generated through the scattering of light neutrinos on the domain wall. However, it may not be the final asymmetry. This is because of the thermally equilibrated lepton violating processes mediated by the right-handed neutrinos can erase the produced asymmetry. Therefore, a final asymmetry and hence the bound on right-handed neutrino masses can only be obtained by solving the Boltzmann equations [18]. We assume a normal mass hierarchy in the right-handed neutrino sector. In this scenario, as the temperature falls, first  $N_3$  and  $N_2$  go out of thermal equilibrium while  $N_1$  is in thermal equilibrium. Therefore, it is the density and mass of  $N_1$  are important in the present case which enter into the Boltzmann equations. The relevant Boltzmann equations for the present purpose are [7]

$$\frac{dY_{N_1}}{dZ} = -(D + S)(Y_{N_1} - Y_{N_1}^{\text{eq}}), \quad (7)$$

$$\frac{dY_{B-L}}{dZ} = -WY_{B-L}, \quad (8)$$

where  $Y_{N_1}$  is the density of  $N_1$  in a comoving volume and that of  $Y_{B-L}$  is the  $B-L$  asymmetry. The parameter  $Z = M_1/T$ . The various terms  $D$ ,  $S$  and  $W$  are representing the decay, scatterings and the wash out processes involving the right-handed neutrinos. In particular,  $D = \Gamma_D/ZH$ , with

$$\Gamma_D = \frac{1}{16\pi v^2} \tilde{m}_1 M_1^2, \quad (9)$$

where  $\tilde{m}_1 = (m_D^\dagger m_D)_{11}/M_1$  is called the effective neutrino mass parameter. Similarly  $S = \Gamma_S/HZ$  and  $W = \Gamma_W/HZ$ . Here  $\Gamma_S$  and  $\Gamma_W$  receives the contribution from  $\Delta_L = 1$  and  $\Delta_L = 2$  lepton violating scattering processes. The dependence of the scattering rates involved in  $\Delta_L = 1$  lepton violating processes on the parameters  $\tilde{m}_1$  and  $M_1$  is similar to that of the decay rate  $\Gamma_D$ . As the Universe expands these  $\Gamma$ 's compete with the Hubble expansion parameter. Therefore in a comoving volume we have

$$\left(\frac{\gamma_D}{sH(M_1)}\right), \left(\frac{\gamma_{\phi,s}^{N_1}}{sH(M_1)}\right), \left(\frac{\gamma_{\phi,t}^{N_1}}{sH(M_1)}\right) \propto k_1 \tilde{m}_1. \quad (10)$$

On the other hand, the dependence of the  $\gamma$ 's in  $\Delta_L = 2$  lepton number violating processes on  $\tilde{m}_1$  and  $M_1$  are given by

$$\left(\frac{\gamma_{N_1}^l}{sH(M_1)}\right), \left(\frac{\gamma_{N_1,t}^l}{sH(M_1)}\right) \propto k_2 \tilde{m}_1^2 M_1. \quad (11)$$

Finally there are also lepton conserving processes where the dependence is given by

$$\left(\frac{\gamma_{Z'}}{sH(M_1)}\right) \propto k_3 M_1^{-1}. \quad (12)$$

In Eqs. (10)–(12),  $k_i$ ,  $i = 1, 2, 3$ , are dimensionful constants determined from other parameters. Since the lepton conserving processes are inversely proportional to the mass scale of  $N_1$ , they rapidly bring the species  $N_1$  into thermal equilibrium for  $T \gg M_1$ . Further smaller the values of  $M_1$ , the washout effects (11) are negligible because of their linear dependence on  $M_1$ . This is the regime in which we are while solving the Boltzmann equations in the following.

Eqs. (7) and (8) are solved numerically. The initial  $B-L$  asymmetry is the net raw asymmetry produced during the  $B-L$  symmetry breaking phase transition by any thermal or non-thermal processes. As such we impose the following initial conditions

$$Y_{N_1}^{\text{in}} = Y_{N_1}^{\text{eq}} \quad \text{and} \quad Y_{B-L}^{\text{in}} = \eta_{B-L}^{\text{raw}}, \quad (13)$$

by assuming that there are no other processes creating lepton asymmetry below the  $B-L$  symmetry breaking scale. This requires  $\Gamma_D \leq H$  at an epoch  $T \geq M_1$  and hence lead to a bound [19]

$$m_\nu < m_* \equiv 4\pi g_*^{1/2} \frac{G_N^{1/2}}{\sqrt{2}G_F} = 6.5 \times 10^{-4} \text{ eV}. \quad (14)$$

Alternatively in terms of Yukawa couplings this bound reads

$$h_\nu \leq 10x, \quad \text{with } x = (M_1/M_{\text{Pl}})^{1/2}. \quad (15)$$

At any temperature  $T \geq M_1$ , wash out processes involving  $N_1$  are kept under check due to the  $\tilde{m}_1^2$  dependence in (11) for small values of  $\tilde{m}_1$ . As a result a given raw asymmetry suffers limited erasure. As the temperature falls below the mass scale of  $N_1$  the wash out processes become negligible leaving behind a final lepton asymmetry. Fig. 1 shows the result of solving the Boltzmann equations for different values of  $M_1$ . An important conclusion comes from this figure is that for smaller values of  $M_1$  the wash out effects are tiny. Hence by demanding the initial raw asymmetry is the required asymmetry of the present Universe we can conspire the mass scale of  $N_1$  to be as low as 1 TeV. For this value of  $M_1$ , using Eq. (15), we get the constraint  $h_\nu \leq 10^{-7}$ . Further lowering of  $M_1$  needs  $h_\nu < 10^{-7}$ . This was the prediction of the model  $\nu\text{MSM}$  to keep  $M_1$  in KeV range, albeit the leptogenesis mechanism was different. However in the present case as we see in Section 4, the bound on  $M_1$  is very much tight from the flavor changing neutral current unless we allow sufficiently small Yukawa couplings.

Note that in Eq. (8) we assume that there are no other sources producing lepton asymmetry below the  $B-L$  symmetry breaking phase transition. This can be justified by considering small values of  $h_\nu$ , since the  $CP$  asymmetry parameter  $\epsilon_1$  depends quadratically on  $h_\nu$ . Hence for  $h_\nu \leq 10^{-7}$  the lepton asymmetry  $Y_L \leq O(10^{-14})$ , which is far less than the raw asymmetry produced by the scatterings of neutrinos on the domain walls. This explains the absence of lepton asymmetry generating term in Eq. (8).

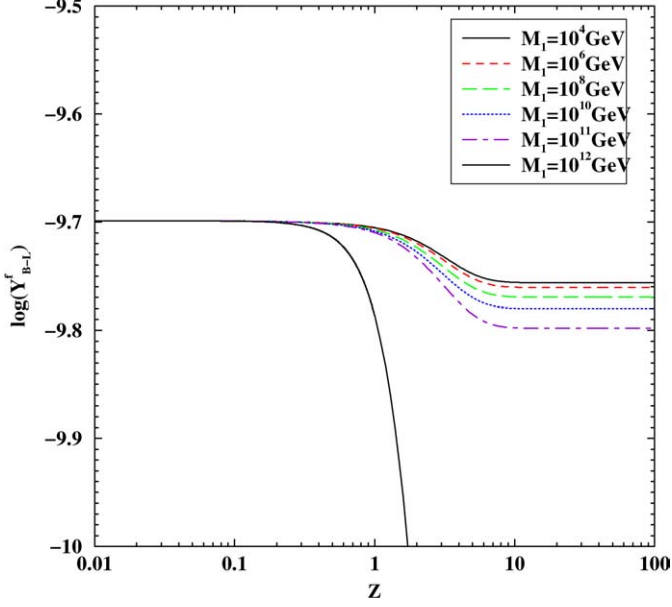


Fig. 1. The evolution of  $B-L$  asymmetry for different values of  $M_1$  shown against  $Z$  ( $= M_1/T$ ) for  $\tilde{m}_1 = 10^{-4}$  eV and  $\eta_{B-L}^{\text{raw}} = 2.0 \times 10^{-10}$ .

### 3. Constraint on the $B-L$ breaking scale

Below the mass scale of  $N_1$  the lepton conserving processes  $N_1 N_1 \rightarrow f \bar{f}$  mediated by the  $Z'$  boson fall out of equilibrium. Here  $f$  and  $\bar{f}$  are the SM fermions and antifermions. The cross-section is given as

$$\sigma\left(N_1 N_1 \rightarrow \sum_f f \bar{f}\right) \sim \frac{1}{4\pi} \frac{E^2}{v_{B-L}^4}, \quad (16)$$

where we have used the mass of  $Z'$  boson  $M_{Z'} = g' v_{B-L}$ , with  $v_{B-L}$  is the  $B-L$  symmetry breaking scale. At the epoch  $T \sim M_1$  the rate of lepton conserving process mediated by the  $Z'$  boson is given by

$$\Gamma_{Z'} = n_{N_1} \langle \sigma v \rangle, \quad (17)$$

where  $n_{N_1}$  is the density of  $N_1$  at that epoch. Further at the epoch  $T \gtrsim M_1$ ,  $n_{N_1} = n_{N_1}^{\text{eq}} = 2T^3/\pi^2$ . Hence substituting it in Eq. (17) and using  $\sigma$  from Eq. (16) we get

$$\Gamma_{Z'} = \frac{1}{2\pi^3} \frac{M_1^5}{v_{B-L}^4}. \quad (18)$$

Requiring  $\Gamma_{Z'} \leq H(M_1)$  we get

$$\begin{aligned} v_{B-L} &\geq \left( \frac{M_{\text{Pl}}}{2\pi^3 \times 1.67 g_*^{1/2}} \right)^{1/4} M_1^{3/4} \\ &\sim 10^6 \text{ GeV} \left( \frac{M_1}{100 \text{ GeV}} \right)^{3/4}. \end{aligned} \quad (19)$$

This tells us that for  $M_1 = 100$  GeV, the  $B-L$  breaking scale is greater than  $10^6$  GeV. This is in well agreement with Eq. (15) for  $h_\nu \leq 10^{-7}$ .

### 4. FCNC constraint on the mass scale of $N_1$

In a flavor basis the Lagrangian describing the neutral current for one generation of fermions is given as

$$\mathcal{L} \simeq \frac{g}{2 \cos \theta_W} Z^\mu \bar{\nu}_{eL} \gamma_\mu \nu_{eL}, \quad (20)$$

where  $\theta_W$  is the weak mixing angle. Rewriting Eq. (20) in a mass basis we get

$$\begin{aligned} \mathcal{L} \simeq \frac{g}{2 \cos \theta_W} Z^\mu & \left[ \cos^2 \theta \bar{\nu}_1 \gamma_\mu L \nu_1 + \sin^2 \theta \bar{N}_1 \gamma_\mu L N_1 \right. \\ & \left. + \cos \theta \sin \theta (\bar{N}_1 \gamma_\mu L \nu + \bar{\nu}_1 \gamma_\mu L N_1) \right], \end{aligned} \quad (21)$$

where  $L$  is the left-handed projection operator and  $\theta$  is the mixing angle and is given by

$$\theta = \frac{m_D}{M_1} = \left( \frac{m_\nu}{M_1} \right)^{1/2}, \quad (22)$$

where we have used Eq. (2). Thus there is a flavor changing neutral current in the model as given by the third term in Eq. (21). This is unlike the case in SM. Hence by requiring  $\theta$  to be small, the flavor changing neutral current can be suppressed. Using the current bound  $m_{\nu_e} \leq 0.6$  eV from the neutrino less double beta decay experiment [20] we get from Eq. (22) that  $\theta \lesssim 10^{-6}$  for  $M_1 \gtrsim 1$  TeV.

On the other hand, if we relax the upper bound on  $\theta$  by three orders larger than the above bound then we get a lower bound on  $M_1$  to be  $\gtrsim 1$  GeV. This will allow the following decay width  $\Gamma(Z \rightarrow \nu N) \sim \theta^2 165 \text{ MeV}$  [21], for mass of  $N_1$  ranging from 1 to 80 GeV. If  $\theta$  is large this decay has a distinctive signature through the decay modes of  $N_1$ . In particular,  $\Gamma(N_1 \rightarrow 3\nu) \propto \theta^2$ . Therefore, the above decay mode of  $Z$  boson is highly restricted.

Now we study the bound on  $\theta$  by considering the magnitude of Dirac Yukawa coupling of the neutrinos. Since  $\theta = m_D/M_1 = h_\nu v/M_1$ , we can achieve small values of  $\theta$  by demanding  $h_\nu \ll h_e$  even for small values of  $M_1$ . This was the prediction of  $\nu$ MSM model, where the Yukawa coupling  $h_\nu$  was required to be very small.

### 5. Dark matter constraint on mass scale of $N_1$

One of the important questions in cosmology is that how much the masses of the galaxies contribute to the critical density

$$\rho_c = \frac{3H_0^2}{8\pi G_N} \equiv 10^4 h^2 \text{ eV/cm}^3 \quad (23)$$

of the present Universe. Here  $H_0 = h \times 100 \text{ km s}^{-1} \text{ Mpc}^{-1}$ , with  $0.4 \lesssim h \lesssim 1.0$  is the Hubble expansion parameter that is observed today. The best fit value, combinly given by the WMAP, 2dFGRS and Lyman  $\alpha$  forest data, is  $h = 0.72 \pm 0.03$  [5]. On the other hand, the  $\Omega$  parameter defined for the total density of the Universe is given by

$$\Omega_{\text{tot}} = \Omega_m + \Omega_\Lambda = 1.02 \pm 0.02, \quad (24)$$

where the various  $\Omega$ 's are defined as  $\Omega_i = (\rho_i/\rho_c)$ . Eq. (24) indicates that the present Universe is flat with the mass density contributed by the galaxies is approximately equal to it's critical density. The best fit value for the present matter component of the Universe, combinely given by the WMAP with 2dFGRS and Lyman alpha forest data, is  $\Omega_m = 0.133 \pm 0.006/h^2$ . However, the baryonic component of matter is found to be  $\Omega_B = 0.0226 \pm 0.0008/h^2$ . This implies that the present Universe contains significant amount of nonbaryonic matter which is given by  $\Omega_{NB} = 0.1104/h^2$ . The missing matters are usually treated as dark matter (DM).

An important issue of the particle physics and cosmology is that the nature of dark matter and its role in the evolution of the Universe. Had it been the cold dark matter it had played an important role in the formation of large scale structure of the Universe. At present the contribution of light neutrinos having masses varying from  $5 \times 10^{-4}$  eV to 1 MeV is [22]

$$\Omega_\nu \leq 0.0076/h^2, \quad 95\% \text{ C.L.} \quad (25)$$

However, this is not sufficient to explain the nonbaryonic component of matter. In the present model we propose that the lightest right-handed neutrino can be a suitable candidate for cold DM for which the life time of  $N_1$  must satisfy the constraint,  $\tau_{N_1} > 2t_0$ , where  $t_0$  is the present age of the Universe. Alternatively we require  $\Gamma_{N_1} < H_0$ , the present Hubble expansion parameter. This gives the constraint on the Dirac mass of the neutrino to be

$$(m_D^\dagger m_D)_{11} < 1.19 \times 10^{-40} \text{ GeV}^2 \left( \frac{10^3 \text{ GeV}}{M_1} \right). \quad (26)$$

A similar constraint on the Dirac mass of the neutrino was obtained in Ref. [23] for  $N_1$  to be a candidate of cold DM.

Since the massive neutrinos are stable in the cosmological time scale we have to make sure that it should not over-close the Universe. For this we have to calculate the density of the heavy neutrino  $N_1$  at the present epoch of temperature  $T_0 = 2.75$  K. The number density of  $N_1$  at present is given by

$$n_{N_1}(T_0) = n_{N_1}(T_D) \left( \frac{T_0}{T_D} \right)^3, \quad (27)$$

where  $T_D$  is the temperature of the thermal bath when the massive neutrinos got decoupled. This can be calculated by considering the out of equilibrium of the annihilation rate  $\Gamma_{\text{ann}}$  of the process  $N_1 N_1 \rightarrow f \bar{f}$ . We assume that at a temperature  $T_D$

$$\Gamma_{\text{ann}}/H(T_D) \simeq 1, \quad (28)$$

where  $\Gamma_{\text{ann}}$  is essentially given by Eq. (17) and

$$H(T_D) = 1.67 g_*^{1/2} \frac{T_D^2}{M_{\text{Pl}}} \quad (29)$$

is the Hubble expansion parameter during the decoupled era. Considering the effective four-Fermi interaction of the annihilation processes  $\sigma$  can be parameterized as [23]

$$\sigma_{N_1} = \frac{G_F^2 M_1^2}{2\pi} c, \quad (30)$$

where  $c$  is the compensation factor and is taken to be  $O(10^{-2})$ . Further  $n_{N_1}$  is the density of  $N_1$  at an epoch  $T \sim M_1$ . At any temperature  $T$ , the density distribution  $n_{N_1}$  is given by

$$n_{N_1}(T) = 2 \left( \frac{M_1 T}{2\pi} \right)^{3/2} \exp\left(-\frac{M_1}{T}\right). \quad (31)$$

Using (29), (30) and (31) in Eq. (28) we get

$$\frac{\Gamma_{\text{ann}}}{H(T_D)} = 1.2 \times 10^{-2} g_*^{-1/2} N_{\text{ann}} G_F^2 M_1^3 M_{\text{Pl}} c z_D^{1/2} \exp(-z_D) \simeq 1, \quad (32)$$

where  $z_D = M_1/T_D$  and  $N_{\text{ann}}$  is the number of annihilation channels which we take  $\approx 10$ . Solving for  $z_D$  from Eq. (32) we get

$$z_D \approx \ln \left[ \frac{N_{\text{ann}}}{82 g_*^{1/2}} (G_F^2 c M_1^3 M_{\text{Pl}}) \right]. \quad (33)$$

Using (32) in Eq. (27) we get

$$\begin{aligned} n_{N_1}(T_0) &= \frac{2}{(2\pi)^{3/2}} z_D^{3/2} \exp(-z_D) T_0^3 \\ &= \frac{2.016 \times 10^{-11}}{\text{cm}^3} \left( \frac{\text{TeV}}{M_1} \right)^3 \\ &\quad \times \left[ 1 + 0.02 + 0.21 \ln \frac{M_1}{1 \text{ TeV}} \right]. \end{aligned} \quad (34)$$

Now we can define the energy density of  $N_1$  at the present epoch as

$$\rho_{N_1} = n_{N_1} M_1 = \frac{20.16}{\text{cm}^3} \left( \frac{1 \text{ TeV}}{M_1} \right)^2 (1 + \text{correction}). \quad (35)$$

Using Eqs. (23) and (35) we can get the  $\Omega$  parameter for  $N_1$  as

$$\Omega_{N_1} = \frac{\rho_{N_1}}{\rho_c} = \left( \frac{0.2016 \times 10^{-2}}{h^2} \right) \left( \frac{1 \text{ TeV}}{M_1} \right)^2. \quad (36)$$

Thus Eq. (36) shows that for  $M_1 = 1$  TeV the contribution of  $N_1$  to the present DM,  $\Omega_{\text{DM}} = (0.1104/h^2)$  is two orders less. On the other hand, if we allow  $M_1 \simeq 100$  GeV [24] then we can satisfy the present DM constraint  $\Omega_{\text{DM}} = (0.1104/h^2)$ . In this mass limit of  $N_1$  we get from Eq. (22) that the mixing angle  $\theta \simeq 10^{-5}$ .

## 6. Summary and conclusion

We studied the dark matter and leptogenesis constraints on the mass scale of lightest right-handed neutrino in a gauged  $B-L$  symmetric model. In this model the break down of the  $B-L$  gauge symmetry produces a net raw lepton asymmetry which under goes a limited erasure for  $\tilde{m}_1 \lesssim 10^{-4}$  eV and  $M_1 < 10^{12}$  GeV and hence leaves behind the required lepton asymmetry which gets converted to the baryon asymmetry that is observed today. Therefore the assumption of raw lepton asymmetry of  $\sim O(10^{-10})$  allows the mass scale of lightest right-handed neutrino to be 1 TeV or less. However, for  $M_1 = 1$  TeV the contribution of  $N_1$  to wards cold DM is two orders less than the required value. On the other hand by requiring the mass scale of lightest right-handed neutrino to be

$O(10^2)$  GeV we can satisfy both leptogenesis as well as cold DM constraint. Further in the left–right symmetric model for  $M_1 = 100$  GeV the mixing angle  $\theta \lesssim 10^{-5}$  and hence the flavor changing neutral current is suppressed.

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