Gauged $B-L$ symmetry and baryogenesis via leptogenesis at TeV scale

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Abstract

It is shown that the requirement of preservation of baryon asymmetry does not rule out a scale for leptogenesis as low as 10 TeV. The conclusions are compatible with see-saw mechanism if for example the pivot mass scale for neutrinos is $\approx 10^{-2}$ that of the charged leptons. We explore the parameter space $\tilde{m}_1$- M_1 of relevant light and heavy neutrino masses by solving Boltzmann equations. A viable scenario for obtaining baryogenesis in this way is presented in the context of gauged $B-L$ symmetry.

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I. INTRODUCTION

It has long been recognized that existence of heavy Majorana neutrinos has important consequences for the baryon asymmetry of the Universe. With the discovery of the neutrino masses and mixings, it becomes clear that only $B-L$ can be considered to be a conserved global symmetry of the Standard Model (SM) and not the individual quantum numbers $B-L_e$, $B-L_\mu$ and $B-L_\tau$. Combined with the fact that the classical $B+L$ symmetry is anomalous [1, 2, 3] it becomes important to analyse the consequences of any $B-L$ violating interactions because the two effects combined can result in the unwelcome conclusion of the net baryon number of the Universe being zero.

At present two broad possibilities exist as viable explanations of the observed baryon asymmetry of the Universe. One is the baryogenesis through leptogenesis [4]. This has been analysed extensively in [5, 6, 7, 8] and has provided very robust conclusions for neutrino physics. Its typical scale of operation has to be high, at least an intermediate scale of 10^9 GeV. This has to do with the intrinsic competition needed between the lepton number violating processes and the expansion scale of the Universe. While the mechanism does not prefer any particular unification scheme, it has the virtue of working practically unchanged upon inclusion of supersymmetry [9].

The alternative to this is provided by mechanisms which work at the TeV scale [10, 11, 12, 13] and rely on the new particle content implied in supersymmetric extensions of the SM. The Minimal Supersymmetric SM (MSSM) holds only marginal possibilities for baryogenesis. The Next to Minimal or NMSSM possesses robust mechanism for baryogenesis [14] however the model has unresolved issues vis a vis the μ problem due to domain walls [15]. However its restricted version, the nMSSM is reported [16, 17, 18, 19, 20] to tackle all of the concerned issues.

It is worth investigating other possibilities, whether or not supersymmetry is essential to the mechanism. In this paper we study afresh the consequence of heavy Majorana neutrinos given the current knowledge of light neutrinos. The starting point is the observation [21, 22] that the heavy neutrinos participate in the erasure of any pre-existing asymmetry through scattering as well as decay and inverse decay processes. Estimates using general behavior of the thermal rates lead to a conclusion that there is an upper bound on the temperature $T_{\mathrm{B-L}}$ at which $B-L$ asymmetry could have been created. This bound is $T_{B-L} \lesssim 10^{13}\text{GeV}\times(1eV/m_\nu)^2$, where m_ν is the typical light neutrino mass. This bound is too weak to be of accelerator physics interest. We extend this analysis by numerical solution of the Boltzmann equations and obtain regions of viability in the parameter space spanned by $\tilde{m}_1$-M_1, where $\tilde{m}_1$ is a light neutrino mass parameter defined in eq. (4) and M_1 is the mass of the lightest of the heavy Majorana neutrinos. We find that our results are in consonance with [22] where it was argued that scattering processes provide a weaker constraint than the decay processes. If the scatterings become the main source of erasure of the primordial asymmetry then the constraint can be interpreted to imply $T_{B-L} < M_1$. Further, this temperature can be as low as TeV range with $\tilde{m}_1$ within the range expected from neutrino observations. This is compatible with see-saw mechanism if the "pivot" mass scale is that of the charged leptons.

In [23, 24] it was shown that the Left-Right symmetric model [25, 26] presents just such a possibility. In this model $B-L$ appears as a gauged symmetry in a natural way. The phase transition is rendered first order so long as there is an approximate discrete symmetry $L \leftrightarrow R$, independent of details of other parameters. Spontaneously generated CP phases then allow creation of lepton asymmetry. We check this scenario here against our numerical results and in the light of the discussion above.

In the following we first recapitulate the arguments concerning erasure of $B-L$ asymmetry due to heavy Majorana neutrinos, then present the relevant numerical results, then discuss the possibility for the Left-Right symmetric model and present a summary in conclusion.

II. ERASURE CONSTRAINTS - SIMPLE VERSION

The presence of several heavy Majorana neutrino species (N_i) gives rise to processes depleting the existing lepton asymmetry in two ways. They are (i) scattering processes (S) among the SM fermions and (ii) Decay (D) and inverse decays (ID) of the heavy neutrinos. We assume a normal hierarchy among the right handed neutrinos such that only the lightest of the right handed neutrinos (N_1) makes a significant contribution to the above mentioned processes. At first we use a simpler picture, though the numerical results to follow are based on the complete formalism. The dominant contributions to the two types of processes are governed by the temperature dependent rates

$$\Gamma_D \sim \frac{h^2 M_1^2}{16\pi(4T^2+M_1^2)^{1/2}} \qquad \text{and} \qquad \Gamma_S \sim \frac{h^4}{13\pi^3}\frac{T^3}{(9T^2+M_1^2)}, \tag{1}$$

where h is typical Dirac Yukawa coupling of the neutrino.

Let us first consider the case $M_1 > T_{B-L}$. For $T < T_{B-L}$, the N_1 states are not populated, nor is there sufficient free energy to create them, rendering the D-ID processes unimportant. We work in the scenario where the sphalerons are in equilibrium, maintaining rough equality of B and L numbers. As the $B-L$ continues to be diluted we estimate the net baryon asymmetry produced as [24]

$$10^{-d_B} \equiv \exp\left(-\int_{t_{B-L}}^{t_{EW}} \Gamma_S dt\right) = \exp\left(-\int_{T_{EW}}^{T_{B-L}} \frac{\Gamma_S}{H}\frac{dT}{T}\right), \tag{2}$$

where t_{B-L} is the time of the $(B-L)$-breaking phase transition, H is the Hubble parameter, and t_{EW} and T_{EW} corresponds to the electroweak scale after which the sphalerons are ineffective. Evaluating the integral gives an estimate for the exponent as

$$d_B \cong \frac{3\sqrt{10}}{13\pi^4 \ln 10\sqrt{g_*}}\, h^4 \frac{M_{Pl}T_{B-L}}{M_1^2}. \tag{3}$$

The same result upto a numerical factor is obtained in [27], the suppression factor $\omega^{(2)}$ therein. Eq. (3) can be solved for the Yukawa coupling h which gives the Dirac mass term for the neutrino: $h^4 \lesssim 3200\, d_B \left(\frac{M_1^2}{T_{B-L}M_{Pl}}\right)$ where we have taken $g_* = 110$ for definiteness and d_B here stands for total depletion including from subdominant channels. This can be further transformed into an upper limit on the light neutrino masses using the canonical seesaw relation. Including the effect of generation mixing, the parameter which appears in the thermal rates is

$$\tilde{m_1} \equiv \frac{(m_D^\dagger m_D)_{11}}{M_1} \tag{4}$$

and is called the *effective neutrino mass* [7]. The constraint (3) can then be recast as

$$m_\nu \ \sim\ \tilde{m}_1 \lesssim \frac{180v^2}{\sqrt{T_{B-L}M_{Pl}}}\left(\frac{d_B}{10}\right)^{1/2} \tag{5}$$

This bound is useful for the case of large suppression. Consider $d_B = 10$. If we seek $T_{B-L} \sim 1$TeV and $M_1 \sim 10$TeV, this bound is saturated for $m_\nu \approx 50$keV, corresponding to $h \approx m_\tau/v$. This bound is academic in view of the WMAP bound $\sum m_{\nu_i} \approx 0.69eV$ [28] . On the other hand, for the phenomenologically interesting case $m_\nu \approx 10^{-2}$eV, with

$h \approx 10^{-5} \approx m_e/v$ and with M_1 and T_{B-L} as above, eq. (5) can be read to imply that in fact d_B is vanishingly small. This in turn demands, in view of the low scale we are seeking, a non-thermal mechanism for producing lepton asymmetry naturally in the range 10^{-10}. Such a mechanism is discussed in sec. IV.

In the opposite regime $M_1 < T_{B-L}$, both of the above types of processes could freely occur. The condition that complete erasure is prevented requires that the above processes are slower than the expansion scale of the Universe for all $T > M_1$. It turns out to be sufficient [22] to require $\Gamma_D < H$ which also ensures that $\Gamma_S < H$. This leads to the requirement

$$m_\nu < m_* \equiv 4\pi g_*^{1/2} \frac{G_N^{1/2}}{\sqrt{2}G_F} = 6.5 \times 10^{-4} eV \tag{6}$$

where the parameter m_* [22] contains only universal couplings and g_*, and may be called the *cosmological neutrino mass.*

The constraint of equation (6) is compatible with models of neutrino mass if we identify the neutrino Dirac mass matrix m_D as that of charged leptons. For a texture of m_D [29, 30]

$$m_D = \frac{m_\tau}{1.054618} \begin{pmatrix} 0 & a & 0 \\ a & 0 & b \\ 0 & b & c \end{pmatrix}, \tag{7}$$

the hierarchy for charged leptons can be generated if

$$a = 0.004, \quad b = 0.24 \quad and \quad c = 1. \tag{8}$$

For these set of values of a, b, and c the above mass matrix m_D is normalized with respect to the τ-lepton mass. Using the mass matrix (7) in equation (4) we get

$$\tilde{m_1} = 4.16 \times 10^{-4} eV \left(\frac{10^8 GeV}{M_1} \right). \tag{9}$$

Thus with this texture of masses, eq. (6) is satisfied for $M_1 \gtrsim 10^8 GeV$. If we seek M_1 mass within the TeV range, this formula suggests that the texture for the neutrinos should have the Dirac mass scale smaller by 10^{-2} relative to the charged leptons.

The bound (6) is meant to ensure that depletion effects remain unimportant and is rather strong. A more detailed estimate of the permitted values of $\tilde{m}_1$ and M_1 is obtained by solving the relevant Boltzmann equations. In the rest of the paper we investigate these for the case $T_{B-L} > M_1$ and $\Gamma_D < H$.

III. SOLUTION OF BOLTZMANN EQUATIONS

The relevant Boltzmann equations for our purpose are [5, 7, 8]

$$\frac{dY_{N1}}{dZ} = -(D+S)\,(Y_{N1} - Y_{N1}^{eq}) \tag{10}$$

$$\frac{dY_{B-L}}{dZ} = -WY_{B-L}, \tag{11}$$

where $Y_i = (n_i/s)$ is the density of the species i in a comoving volume, $Z = M_1/T$ and $s = 43.86(g_*/100)T^3$ is the entropy density at an epoch of temperature T. The two terms D and S on the right hand side of equation (10) change the density of N_1 in a comoving volume while the right hand side of equation (11) accounts for the wash out or dilution effects. We assume no new processes below T_{B-L} which can create lepton asymmetry. In a recent work [31] it has been reported that the thermal corrections to the above processes as well as the processes involving the gauge bosons are important for final L-asymmetry. However their importance is under debate [32]. In this paper we limit ourselves to the same formalism as in [7, 8]. In equation (10) $D = \Gamma_D/HZ$, where Γ_D determines the decay rate of N_1, $S = \Gamma_S/HZ$, where Γ_S determines the rate of $\Delta_{\rm L} = 1$ lepton violating scatterings. In equation (11) $W = \Gamma_W/HZ$, where Γ_W incorporates the rate of depletion of the B-L number involving the lepton violating processes with $\Delta_{\rm L} = 1$, $\Delta_{\rm L} = 2$ as well as inverse decays creating N_1. The various Γ's are related to the scattering densities [5] γs as

$$\Gamma_i^X(Z) = \frac{\gamma_i(Z)}{n_X^{eq}} \tag{12}$$

The dependence of the scattering rates involved in $\Delta_{\rm L} = 1$ lepton violating processes on the parameters $\tilde{m}_1$ and M_1 is similar to that of the decay rate Γ_D. As the Universe expands these Γ's compete with the Hubble expansion parameter. Therefore in a comoving volume we have

$$\left(\frac{\gamma_D}{sH(M_1)}\right), \left(\frac{\gamma_{\phi,s}^{N1}}{sH(M_1)}\right), \left(\frac{\gamma_{\phi,t}^{N1}}{sH(M_1)}\right) \propto k_1 \tilde{m}_1. \tag{13}$$

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

$$\left(\frac{\gamma_{N1}^{l}}{sH(M_1)}\right), \left(\frac{\gamma_{N1,t}^{l}}{sH(M_1)}\right) \propto k_2 \tilde{m}_1^2 M_1. \tag{14}$$

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

$$\left(\frac{\gamma_{Z'}}{sH(M_1)}\right) \propto k_3 M_1^{-1}. \tag{15}$$

In the above equations (13), (14), (15), 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 $M_1 < T$. Further, for the smaller values of M_1, the washout effects (14) 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.

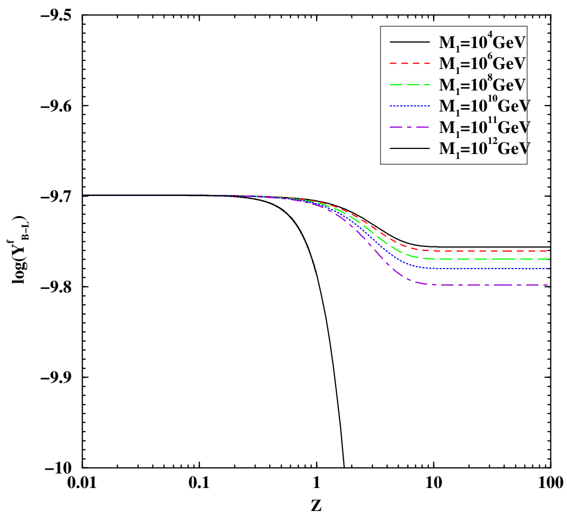

FIG. 1: The evolution of B-L asymmetry for different values of M_1 shown against $Z(= M_1/T)$ for $\tilde{m}_1 = 10^{-4}$eV and $\eta^{raw} = 2.0 \times 10^{-10}$

The equations (10) and (11) 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 process. As such we impose the following initial conditions

$$Y_{N1}^{in} = Y_{N1}^{eq} \;\; and \;\; Y_{B-L}^{in} = \eta_{B-L}^{raw}. \tag{16}$$

At temperature $T \geq M_1$, wash out effects involving N_1 are kept under check due to the $\tilde{m}_1^2$ dependence in (14) 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 M_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.

If we demand that the initial raw asymmetry is of the order of $\alpha \times 10^{-10}$, with $\alpha \sim O(1)$, then in order to preserve the final asymmetry of the same order as the initial one it is

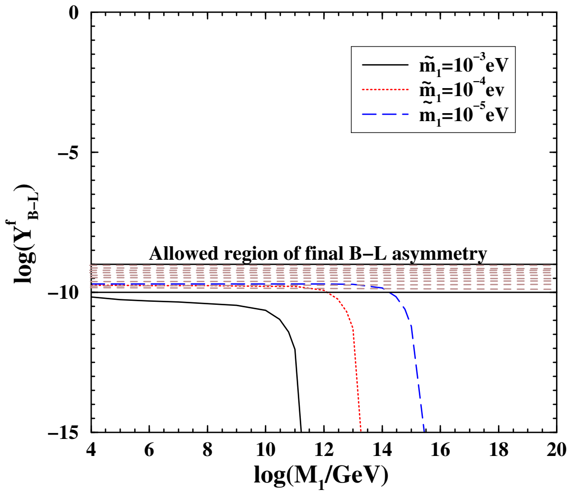

FIG. 2: The allowed values of M_1 against the required final asymmetry is shown for $\eta^{raw} = 2.0 \times 10^{-10}$

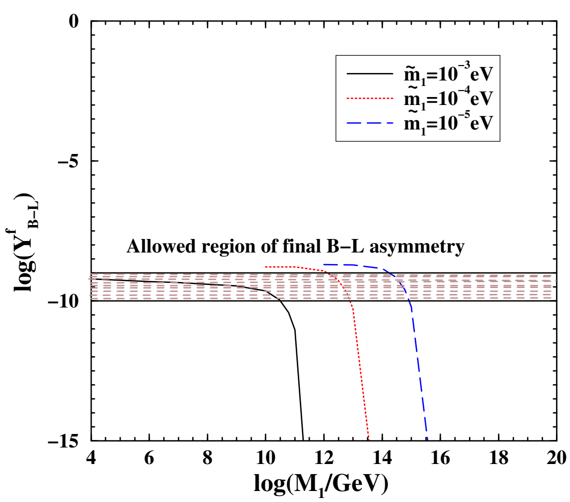

FIG. 3: The allowed values of M_1 against the final required asymmetry is shown for $\eta^{raw} = 2.0 \times 10^{-9}$

necessary that the neutrino mass parameter $\tilde{m}_1$ should be less than $10^{-3}eV$. This can be seen from fig. 2. For $\tilde{m}_1 = 10^{-3}eV$ we can not find any value of M_1 to preserve the final asymmetry, $\alpha \times 10^{-10}$ in the allowed region. This is because of the large wash out effects as inferred from the equation (14). However, for $\tilde{m}_1 = 10^{-4}eV$ we get a *lowest threshold* on the lightest right handed neutrino of the order $10^{12}GeV$. For any value of $M_1 \leq 10^{12}GeV$ the final asymmetry lies in the allowed region. This bound increases by two order of magnitude

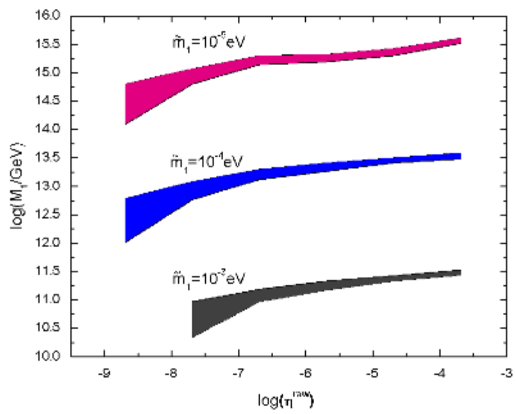

FIG. 4: The allowed region of M_1 is shown for different values of $\tilde{m}_1$ for large values of η^{raw}

for further one order suppression of the neutrino parameter $\tilde{m}_1$. The important point being that $M_1 = 10TeV$ is within the acceptable range.

We now consider the raw asymmetry one order more than the previous case *i.e.* $\eta^{raw}_{B-L} = \alpha \times 10^{-9}$. From fig. 3 we see that for $\tilde{m}_1 = 10^{-3}eV$ there is only an upper bound $M_1 = 10^{10.5}GeV$, such that the final asymmetry lies in the allowed region for all smaller values of M_1. Thus for the case of raw asymmetry an order of magnitude smaller, the upper bound on M_1 decrease by two orders of magnitude (e.g. compare previous paragraph). However, the choice of smaller values of $\tilde{m}_1$ leads to a small window for values of M_1 for which we end up with the final required asymmetry. In particular for $\tilde{m}_1 = 10^{-4}eV$ the allowed range for M_1 is $(10^{12} - 10^{13})GeV$, while for $\tilde{m}_1 = 10^{-5}eV$ the allowed range shifts to $(10^{14} - 10^{15})GeV$. The window effect can be understood as follows. Increasing the value of M_1 tends to lift the suppression imposed by the $\tilde{m}_1^2$ dependence of the wash out effects, thus improving efficiency of the latter. However, further increase in M_1 makes the effects too efficient, erasing the raw asymmetry to insignificant levels.

The windowing effect emerges clearly as we consider the cases of large raw asymmetries. This is shown in fig. 4. It is seen that as the raw asymmetry increases the allowed regions become progressively narrower and lie in the range $(10^{10} - 10^{15})GeV$. Thus a given raw lepton asymmetry determines a corresponding small range of the heavy Majorana neutrino masses for which we can obtain the final asymmetry of the required order $\alpha \times 10^{-10}$. Again smaller is the effective neutrino mass $\tilde{m}_1$ larger is the mean value of the allowed mass of the heavy Majorana neutrino and this is a consequence of normal see-saw.

Finally, in the following, we give an example for non-thermal creation of L-asymmetry in the context of left-right symmetric model.

IV. LEPTON ASYMMETRY IN LEFT-RIGHT SYMMETRIC MODEL

We discuss qualitatively the possibility of lepton asymmetry during the left-right symmetry breaking phase transition [24]. In the following we recapitulate the important aspects of left-right symmetric model for our purpose and the possible non-thermal mechanism of producing raw lepton asymmetry. This asymmetry which gets converted to baryon asymmetry, can be naturally small if the quartic couplings of the theory are small. Smallness of zero-temperature CP phase is not essential for this mechanism to provide small raw L asymmetry.

A. Left-Right symmetric model and transient domain walls

In the left-right symmetric model the right handed chiral leptons, which are singlet under the Standard Model gauge group $SU(2)_L \otimes U(1)_Y$, get a new member ν_R per family under the gauge group $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. The augmented symmetry in the Higgs sector requires two additional triplets Δ_L and Δ_R and a bidoublet Φ, which contains two copies of SM Higgs forming an $SU(2)_R$ representation.

The Higgs potential of the theory naturally entails a vacuum structure wherein at the first stage of symmetry breaking, either one of Δ_L or Δ_R acquires a vacuum expectation value the left-right symmetry, $SU(2)_L \leftrightarrow SU(2)_R$, breaks. It is required that Δ_R acquires a VEV first, resulting in $SU(2)_R \otimes U(1)_{B-L} \to U(1)_Y$. Finally $Q = T_L^3 + T_R^3 + \frac{1}{2}(B-L)$, survives after the bidoublet ϕ and the Δ_L acquire VEVs.

If the left-right symmetry were exact, the first stage of breaking gives rise to stable domain walls [33, 34, 35] interpolating between the L and R-like regions. By L-like we mean regions favored by the observed phenomenology, while in the R-like regions the vacuum expectation value of Δ_R is zero. Unless some non-trivial mechanism prevents this domain structure, the existence of R-like domains would disagree with low energy phenomenology. Furthermore, the domain walls would quickly come to dominate the energy density of the Universe. Thus in this theory a departure from exact symmetry in such a way as to eliminate the R-like regions is essential.

The domain walls formed can be transient if there exists a slight deviation from exact discrete symmetry. As a result the thermal perturbative corrections to the Higgs field free energy will not be symmetric and the domain walls will be unstable. This is possible if the low energy ($\sim 10^4$GeV-10^9GeV) left-right symmetric theory is descended from a Grand Unified Theory (GUT) and the effect is small, suppressed by the GUT mass scale. In the process of cooling the Universe would first pass through the phase transition where this approximate symmetry breaks. The slight difference in free energy between the two types of regions causes a pressure difference across the walls, converting all the R-like regions to L-like regions. Details of this dynamics can be found in ref. [24].

B. Leptogenesis mechanism

At least two of the Higgs expectation values in L-R model are generically complex, thus making it possible to achieve natural i.e., purely spontaneous CP violation [36] permitting all parameters in the Higgs potential to be real. It was shown in [24] that within the thickness of the domain walls the net CP violating phase becomes position dependent. Under these circumstances a formalism exists [37, 38, 39], wherein the chemical potential μ_L created for the Lepton number can be computed as a solution of the diffusion equation

$$-D_\nu \mu_L'' - v_w \mu_L' + \theta(x)\,\Gamma_{\rm hf}\,\mu_L = S(x). \tag{17}$$

Here D_ν is the neutrino diffusion coefficient, v_w is the velocity of the wall, taken to be moving in the $+x$ direction, $\Gamma_{\rm hf}$ is the rate of helicity flipping interactions taking place in front of the wall (hence the step function $\theta(x)$), and S is the source term which contains derivatives of the position dependent complex Dirac mass.

After integration of the above equation and using inputs from the numerical solutions we find the raw lepton asymmetry [24]

$$\eta_L^{\rm raw} \cong 0.01\, v_w \frac{1}{g_*} \frac{M_1^4}{T^5 \Delta_w} \tag{18}$$

where $\eta_L^{\rm raw}$ is the ratio of n_L to the entropy density s. In the right hand side Δ_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 Δ_T is the temperature dependent VEV acquired by the Δ_R in the phase of interest, and $\Delta_w^{-1} = \sqrt{\lambda_{eff}} \Delta_T$ in equation (18) we get

$$\eta^{\rm raw}_{B-L} \cong 10^{-4} v_w \left(\frac{\Delta_T}{T}\right)^5 f_1^4 \sqrt{\lambda_{eff}}. \tag{19}$$

Here we have used $g_* = 110$. Therefore, depending on the various dimensionless couplings, the raw asymmetry can take a range a values of $O(10^{-4} \ - \ 10^{-10})$. In particular, the cases $f_1 < 1$ or $f_1 \ll 1$ implicit in our numerical calculation are compatible with above formula.

V. SUMMARY AND CONCLUSION

In this paper we assume non-thermal production of raw lepton asymmetry during the $B - L$ gauge symmetry breaking phase transitions. If this asymmetry passes without much dilution to be the currently observed baryon asymmetry consistent with WMAP and Big Bang nucleosynthesis, then the effective neutrino mass parameter $\tilde{m}_1$ must be less than $10^{-3} eV$. Solution of the relevant Boltzmann equations shows that for $\tilde{m}_1 = 10^{-4} eV$ the mass of lightest right handed neutrino N_1 has to be smaller than $10^{12} GeV$ and can be as low as 10 TeV. In a more restrictive scenario where the neutrino Dirac mass matrix is identified with that of the charged leptons it is necessary that $M_1 > 10^8 GeV$ in order to satisfy $\tilde{m}_1 < 10^{-3} eV$. Therefore in the more restricted scenario all values M_1, $10^8 GeV < M_1 < 10^{12} GeV$ can successfully create the required asymmetry. If the Dirac mass scale of neutrinos is less restricted, much lower values of M_1 are allowed. In particular, a right handed neutrino as low as 10 TeV is admissible.

If the raw asymmetry is large, the numerical solutions show a small window for M_1 to get the final asymmetry of the required order. The allowed range gets smaller as the raw asymmetry gets larger. This is true for all allowed values of the neutrino mass parameter $\tilde{m}_1$.

In summary, if the $B-L$ gauge symmetry is gauged, we start with a clean slate for $B-L$ number and an asymmetry in it can be generated by a non-perturbative mechanism at the scale where it breaks. The presence of heavy right handed neutrinos still permits sufficient asymmetry to be left over in the form of baryons for a large range of values of the $B - L$ breaking scales. While other mechanisms of leptogenesis become unnatural below 10^8 GeV this mechanism even tolerates TeV scale. A specific mechanism of this kind is possible in the context of Left-Right symmetric model, presumably embedded in the larger unifying group $SO(10)$. Upon incorporation of supersymmetry, the qualitative picture remains unaltered. In the simplest possibility, the scale of SUSY breaking is higher than the $B-L$ breaking scale, in which case the present considerations will carry through without any change. The gravitino bound of 10^9 GeV for reheating temperature after inflation is easily accommodated.

VI. ACKNOWLEDGMENT

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