Gauged $B - L$ Leptogenesis

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We propose a new leptogenesis scenario in a gauged $B - L$ model with supersymmetry at the TeV energy scale. Instead of relying on the very small Yukawa couplings of the singlet neutrinos $N^c$ to generate the observed baryon asymmetry of the Universe, which requires a very large resonance enhancement, their $B - L$ gauge interactions are invoked. Successful leptogenesis is then possible if a particular scalar bilinear $\tilde{N}^c \tilde{N}^c$ term is disallowed.

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The current measurement of the baryon-to-entropy ratio of the Universe is given by

$$\frac{Y_B}{s} = \left(0.87 \pm 0.02\right) \times 10^{-10},$$

where $s = 2\pi^2 g_* T^3/45$ is the entropy density and $g_*$ is the effective number of relativistic degrees of freedom. CP violation is an essential requirement in order to obtain this asymmetry. Leptogenesis [2] is the most promising mechanism to explain it. It is known that there are several scenarios of the leptogenesis [3–6].

Leptogenesis through the decay of a heavy singlet neutrino is considered as the best scenario for understanding the observed baryon asymmetry of the Universe. However, the energy scale involved in a successful application is usually in the range $10^9$ to $10^{13}$ GeV, which renders the idea impossible to verify experimentally. It has also been suggested that this mechanism works just as well at the more easily accessible TeV energy scale, but then the very small Yukawa couplings required by neutrino masses implies that this effect is much too small to be visible, unless it is compensated by a very large resonance enhancement [2], i.e. the near mass degeneracy of two singlet neutrinos. As an alternative solution, instead of the resonance-enhancement hypothesis, we suggest that the source of this matter-antimatter asymmetry is actually a gaugino interaction of gauged $U(1)_{B-L}$ symmetry in a supersymmetric extension of the Standard Model (SM) at the TeV scale.

In supersymmetry, the addition of the singlet superfield $N^c$ with $B - L = 1$ implies a fermion $N^c$ and a scalar $\tilde{N}^c$. As $B - L = 1$ is spontaneously broken by singlet superfields $\tilde{\chi}_{1,2}$ with $B - L = \mp 2$, an exact $Z_2$ residual symmetry remains, i.e. $R$ parity, with $R = (-)^3(B - L + 2)$. As a result, $N^c$ acquires a large Majorana mass through $\langle \chi_1^1 \rangle N^c \tilde{N}^c$, so that it may decay into both leptons and antileptons, thereby initiating leptogenesis. As for $\tilde{N}^c$, there are in general two kinds of mass terms: $(\tilde{N}^c)^* \tilde{N}^c$ and $N^c \tilde{N}^c$. If the latter is absent, then $\tilde{N}^c$ may be assigned $L = -1$ in a subset of its interactions. This will be the key to having a successful leptogenesis scenario, using $B - L$ gauge interactions. Some of us have studied resonant scenarios in TeV scale $B - L$ model in Ref.[5].

Consider the two families $N^c_{1,3}$ and $\tilde{N}^c_{1,3}$ with masses arranged in the order

$$M_{N^c_3} < M_{N^c_1} < M_{N^c_5} < M_{N^c_1'},$$

and with $N^c_1$ coupling to $g_{B-L} \tilde{Z}_{B-L}(\tilde{N}^c_1 \cos \theta + \tilde{N}^c_3 \sin \theta)$ and $N^c_3$ coupling to the orthogonal combination, where $g_{B-L}$ is the $B - L$ gauge coupling, and $\tilde{Z}_{B-L}$ is the $B - L$ gaugino which is also assumed to be lighter than the mass difference between $N^c_1$ and $N^c_3$. The decay of $N^c_1$ is then only into $\tilde{N}^c_3 + \tilde{Z}_{B-L}$, with coupling $g_{B-L} \sin \theta$. Since $\tilde{N}^c_3$ represents the misalignment of the two families after supersymmetry breaking, it may be assumed to be very small, i.e. of order $10^{-6}$, to satisfy the out-of-equilibrium condition for $M_{N^c_1}$ at the TeV scale. A large lepton asymmetry proportional to $(g_{B-L} \cos \theta)^2$ may then be generated through the one-loop exchange of $N^c_3$, provided that below $M_{N^c_1}$, additive lepton number is conserved, i.e. $\tilde{N}^c_3$ having $L = -1$ in all its subsequent interactions. In the following we will show in detail how this all works.

As shown in Ref.[7], after the $B - L$ symmetry breaking by the VEVs $\langle \chi_{1,2} \rangle = v_{1,2}^i \cos \theta + v_{1,2}^j \sin \theta$, a bilinear coupling $B_{N^c_1} \tilde{N}^c_{N^c_3}$ is generally obtained and it is given by $B_{N}^2 = -v_{1,2}^i Y_{1,2}^i + v_{1,2}^j Y_{1,2}^j + m_{N}^2 \cos 2\theta$, and they have lepton numbers $L = \mp 1$ respectively. Moreover, if $\cos 2\theta$ is negative, $\tilde{N}^c_3$ can be lighter than $N^c_3$. Actually, only one $\tilde{N}^c_3$ mass eigenstate needs to have lepton number and be lighter than the lightest $N^c$. This is the crucial assumption of our proposal. For $B - L$ neutrinos $\tilde{\chi}_{a} = (\tilde{\chi}_{1,2}, -i \tilde{Z}_{B-L})$, the mass eigenstates $\tilde{\chi}_{ph,y,a}(a = 1, 2, 3)$ are given by the unitary diagonalization matrix $R$ as $\tilde{\chi}_{a} = \sum_b R_{ab} \tilde{\chi}_{ph,y,b}$. $R^t R = 1$. In our
In the numerical calculation, we derive mass eigenvalues and mixing parameter $R$ in the following two limiting cases: Case A) $µ' \gg MB-L$, Case B) $MB-L, M_{Z_{B,L}} \gg µ'$, where $µ'$, $MB-L$, and $M_{Z_{B,L}}$ are defined as the mass parameter of $\chi_{1,2}$, $Z_{B,L}$, and $Z_{B,L}$, respectively.

The Lagrangian, in flavor eigenstates, relevant for our analysis is given by

$$\mathcal{L} = -\sqrt{2}g_{B-L}(\bar{u}_LNY_i^c(N^c_i) - Y_{Nij}\bar{N}_i(N^c_i)\bar{N}_j^c - M^2_{N_{i\bar{j}}}\bar{N}_i^cN_j^c - \frac{1}{2} M_{Nij}(N^c_i)(N^c_j) + h.c., (3)$$

where

$$M_{Nij} = Y_{Nij}u^\prime sin\theta,$$

$$M^2_{N_{i\bar{j}}} = (M_N^2M_{\Gamma})_{ij} + m^2_{N_{i\bar{j}}} + \frac{1}{4} M^2_{Z_{B,L}} cos2\theta_{ij}.$$

These mass matrices are diagonalized by unitary matrices $U$: $U^\dagger M_{\Gamma}U = diag, \Gamma^2 M_{\Gamma}^2 = diag$, and mass eigenstates $(N^c_i)$ and $(\tilde{N}^c_i)$ are defined as $(N^c_i) = U_{ij}(N^c_j)^m, (\tilde{N}^c_i) = \tilde{\Gamma}_{ij}(N^c_j)^m$. Notice that the mixing matrix $U$ and $\tilde{\Gamma}$ are in general different from each other. Therefore, the combination $U\tilde{\Gamma}$ is not unit matrix, and complex. This is the origin of CP violation.

The Lagrangian in mass eigenstate (hereafter we remove the index "i") is given by

$$\mathcal{L} = -A_{aij}\bar{\Psi}_aP_LN_i\bar{N}_j^c - B_{aij}\bar{\Psi}_aP_LN_i\bar{N}_j^c + h.c.$$

where

$$A_{aij} = \sqrt{2}g_{B-L}R_{ia}(\Gamma^1U)_{ij}, B_{aij} = Y_{Nij}R_{ia}(U^\dagger\Gamma)_{ij}. (7)$$

The four-component Majorana spinors are defined as $\Psi = (\bar{\Phi}_{\nu p}\nu_a, \bar{\phi}_{\nu p}\nu_a)^T$ and $N_i = (N^c_i, \tilde{N}^c_i)^T$. Notice that $A, B = O(1)$ naturally for $i = j$ and small for $i \neq j$ as shown below Eq. (12).

Now, we consider leptonogenesis by $N_1 \rightarrow \Psi \tilde{N}_3^c$ induced by Eq. (6), assuming the mass hierarchy $M_{N1} < M_{N2,3}$. We assume that only the lightest $B-L$ neutrino $\Psi \equiv \Psi$ and sneutrino of the third generation $\tilde{N}_3$ are lighter than $N_1$, and satisfy the relation $m_\nu + M_{\tilde{N}_3} < M_{N1}$. Moreover, we restrict ourselves to two limiting cases. Since $R_{11}(R_{13}) < 1$ in the case A (B), $A_{aij}(B_{aij})$, $(a = 1, i = 1, j = 3)$ gives dominant contributions, and $\Psi$ is nearly $Z_{B-L}(\chi_1)$. As emphasized, due to the fact that $B_{N} \approx 0$, $\tilde{N}_3$ carries lepton number, hence the decay $N_1 \rightarrow \Psi \tilde{N}_3^c$ violates lepton number.

CP asymmetry of $N_1 \rightarrow \Psi \tilde{N}_3^c$ decay processes is generated by the interference between tree and one-loop level diagrams of vertex and self-energy correction shown in FIG. 1. It is defined as

$$\epsilon_1 = \frac{\Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c) - \Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c)}{\Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c) + \Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c)}. (8)$$

FIG. 1: Tree and one-loop decay of $N_1 \rightarrow \Psi \tilde{N}_3^c$ decay.

The decay rate $\Gamma$ at one-loop level is given by

$$\Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c) = \frac{1}{2M_{N1}} |\mathcal{A}_{tree}+\mathcal{A}_{loop}|^2 I_2(N_1 \rightarrow \Psi \tilde{N}_3^c). (9)$$

$$\Gamma(N_1 \rightarrow \Psi \tilde{N}_3^c) = \frac{1}{2M_{N1}} |\mathcal{A}_{tree}+\mathcal{A}_{loop}|^2 I_2(N_1 \rightarrow \Psi \tilde{N}_3^c), (10)$$

where the phase space integral of two-body decay $I_2$ is given by

$$I_2(X \rightarrow Y Z) = \frac{1}{8\pi M_X^2} \sqrt{[M_X^2 - (M_Y + M_Z)^2][M_X^2 - (M_Y - M_Z)^2].}$$

$\mathcal{A}_{tree,loop}$ are tree and loop level amplitudes, and $F$ is kinematical factor. As the loop-level diagrams have vertex and self-energy corrections, we write $\mathcal{A}_{loop}$ F as: $\mathcal{A}_{loop} = A_s F_s + A_s F_s$. From Eqs. (9) and (10), the total CP asymmetry $\epsilon_1 = \epsilon_1 + \epsilon_1^\prime$ is given by

$$\epsilon_{1(s)} = \frac{-2Im[A_{tree}\epsilon_{1(s)}]}{|A_{tree}|^2}. (11)$$

with

$$|A_{tree}|^2 = M_{N1}(1 + r_\chi - r_{\tilde{N}_3}) |A_{113}|^2,$$

$$A_{tree}A_s = \sum_k M_{N1}M_{Nk}(A_{113}A_{1k3})^2,$$

$$F_s = \frac{1}{(4\pi)^2} \int dx dy dz \delta(x + y + z - 1) \times \frac{(y - 1)(1 + r_\chi - r_{\tilde{N}_3}) + 2r_\chi x - y z r_{\tilde{N}_3} - z x r_\chi + x z r_{\tilde{N}_3} + y_r r_\chi + z r_N k)}{1 - r_{\tilde{N}_3}}.$$
1. $N_1^c \to \Psi N_3^c$ decay generates $\tilde{N}_3^c$ asymmetry $Y_{\Delta \tilde{N}_3^c}$.

2. $\tilde{N}_3^c$ decays into (s)lepton by Dirac Yukawa couplings, soft SUSY breaking A-term and $\mu$-term, and resulting (s)lepton asymmetry $Y_{\Delta L} (\Delta L)$ is obtained by solving the Boltzmann equations.

3. Sphaleron converts total lepton asymmetry $Y_L = Y_{\Delta L} + Y_{\Delta \tilde{L}}$ to baryon asymmetry $Y_B$.

Moreover, we take into account scattering processes mediated by $B - L$ gauge boson: $N_1 N_1 \to Z_{B-L} \to f \bar{f}$. For the elastic scattering ($f = N_1$), the scattering rate is large for high temperature $z \ll 1$, which realizes kinetic equilibrium. For very large $M_{Z_{B-L}}$, the scattering is Boltzmann suppressed near $z = 1$. So decay dominates at this temperature, and leptogenesis occurs at $z = 1$. This condition may give lower bound as $M_{Z_{B-L}} \gtrsim 10^2 M_{N_1}$ [9,11]. On the other hand for small $M_{Z_{B-L}}$ case, scattering contributions survive until $z \sim 10$. As a result, since asymmetry due to the decay starts to be produced by small $N_1$ abundance at large $z$, only small lepton asymmetry is created unless CP asymmetry is large. On the other hand, since scattering processes by $B - L$ gaugino: $N_1 \tilde{N}_3^c \to Z_{B-L} \to f \bar{f}$, are well suppressed by small mixing matrix $\Gamma_{13} \ll 1$, we neglect them.

The thermal average decay and scattering rates that contribute to Boltzmann equations, which we solved numerically to get the total lepton asymmetry, are give by:

\[
\gamma_D = n_{N_1} \frac{K_1(z)}{K_2(z)} \left[ \Gamma(N_1 \to \tilde{N}_3^c \Psi) + \Gamma(N_1 \to \tilde{N}_3^c \tilde{\Psi}) \right],
\]

\[
\gamma_{N_1 \to L} = n_{N_1} \frac{K_1(\sqrt{\frac{m_{N_1}}{m_L}})}{K_2(\sqrt{\frac{m_{N_1}}{m_L}})} \Gamma(\tilde{N}_3^c \to \bar{L} \bar{H}_2),
\]

\[
\gamma_{N_3 \to \tilde{L}} = n_{N_3} \frac{K_1(\sqrt{\frac{m_{N_3}}{m_{\tilde{L}}}})}{K_2(\sqrt{\frac{m_{N_3}}{m_{\tilde{L}}}})} \Gamma(\tilde{N}_3^c \to \tilde{L} \tilde{H}_2),
\]

\[
\gamma_S = \langle \sigma \rangle = \frac{T}{64\pi^2} \int_{s_{\text{max}}}^{\infty} ds \sqrt{s} \sigma(s)K_1(\frac{\sqrt{s}}{T}),
\]

where $K_1(z)$ and $K_2(z)$ are modified Bessel functions, and $s_{\text{max}} = \max[4 M_{N_1}^2, 4 m_{\tilde{L}}^2]$. The decay rate of $\tilde{N}_3^c$ at $T = 0$ into $\bar{L} \bar{H}_2, \tilde{L} \tilde{H}_2, \tilde{L} \tilde{H}_1$, and $\bar{L} \bar{H}_2$ is written as

\[
\Gamma(\tilde{N}_3^c \to AB) = \sum_{i} \frac{1}{M_{N_3}^2} |(Y_{AB})_{3i}|^2 I_2(\tilde{N}_3^c \to AB),
\]

where the associated couplings $Y_{AB}$ are given by $(\Gamma^T Y_{\nu})_{3i}(M_{N_3}^2 - m_{\nu}^2)^{1/2}$, $(\Gamma^T A_\nu)_{3i}$, $\mu^* (\Gamma^T Y_{\nu})_{3i}$ and $(\Gamma^T M_N Y_{\nu})_{3i}$. The reduced cross section $\sigma(s)$ for fermionic (bosonic) final states $N_1 N_1 \to Z_{B-L} \to \psi \bar{\psi}(\phi \bar{\phi})$ is given by

\[
\sigma(s) = \frac{g_{B-L}^4}{3\pi} \left( \frac{1}{s - M_{Z_{B-L}}^2} + M_{Z_{B-L}}^2 \Gamma_{Z_{B-L}}^2 \right) \frac{1}{s - 4M_{N_1}^2 + 3m_{\tilde{L}}^2 + 10M_{N_1}^2 m_{\tilde{L}}^2} \right),
\]

where $Q_{\psi,\phi}$ is $U(1)_{B-L}$ charge of the field $\psi$ and $\phi$.

Now we give numerical examples for $\nu' = 6$ TeV, $M_N = (5, 5.5, 6)$ TeV, $M_{N_1} = (5.7, 6.1, 0.3)$ TeV, $g_{B-L} = 1$, $M_{Z_{B-L}} = 2 \sqrt{2} g_{B-L} \nu' \approx 17$ TeV, $g_s = 251.25$. We focus on the above two cases A and B with the following inputs:

(A) : $M_{B-L} = 300$ GeV, $\theta = \pi/2$, $\mu' = 0.9 M_{Z_{B-L}}$,

(B) : $M_{B-L} = 1.2 M_{Z_{B-L}}$, $\theta = \pi/4$, $\mu' = 300$ GeV,

Since $Y_N = M_N/(v' \sin \theta)$ is diagonal, $U = 1$. For both cases, the scalar mass matrix $m_{N_1}^2$ of soft SUSY breaking terms has small deviation from the diagonal form, which gives small $\Gamma_{13}$. In order to obtain light $\tilde{N}_3^c$, $(m_{N_3}^2)_{33}$ is tuned to be $(60 \text{ TeV})^2$. The corresponding CP asymmetry $\epsilon_1$ and the out-of-equilibrium condition $\Gamma/H(z = 1)$ are given by

\[
(A) : \epsilon_1 = -0.10, \quad \frac{\Gamma}{H(z = 1)} = 20.0.
\]

\[
(B) : \epsilon_1 = -0.080, \quad \frac{\Gamma}{H(z = 1)} = 12.4.
\]

FIG. [2] show the behavior of $Y_{N_1}$ and $Y_{\Delta \tilde{N}_3^c, \Delta L, \Delta \tilde{L}, \epsilon_1}$ for the case A and B. Sphaleron processes are in equilibrium above the critical temperature $T_c$. In this region, lepton asymmetry is converted into baryon asymmetry by the rate $Y_B = -8/15Y_L$. Below $T_c$, sphaleron processes are still in equilibrium and the conversion rate from lepton to baryon asymmetry is a function of the temperature-dependent VEV $v(T)$ [13]. At some temperature $T_d < T_c$, sphaleron processes are switched off due to the Boltzmann factor and baryon asymmetry never evolves below $T_d$ while lepton asymmetry still evolves by the Boltzmann equations. However, we make approximation that sphaleron processes are active for $T > 100$ GeV, and switched off for $T < 100$ GeV. From this approximation, we obtain the final results with $Y_L = 3 \times 10^{-8}$:

\[
(A) : Y_B = 3.9 \times 10^{-10},
\]

\[
(B) : Y_B = 1.6 \times 10^{-10}.
\]
Therefore we can obtain enough baryon asymmetry.

In conclusion, we have shown that a successful TeV scale leptogenesis can take place in gauged $B-L$ supersymmetric model. In this model, if the right-sneutrino bilinear term is absent, then the lightest sneutrino is assigned a lepton number. Therefore if $\tilde{N}_2^c$ is lighter than $N_1^c$ and scalar mass matrix of $\tilde{N}_1^c$ is almost diagonal, a large lepton asymmetry can be generated by $B-L$ neutralino interactions of $O(1)$ couplings $g_{B-L}$ and/or $Y_N$ through the one-loop exchange of $N_3^c$ for the decay $N_1^c \rightarrow \tilde{N}_3^c \Psi$. This asymmetry of $\tilde{N}_3^c$ is transmitted into asymmetry of lepton and slepton through the Yukawa coupling, trilinear coupling, and $\mu$-term, and sphaleron converts lepton asymmetry to baryon asymmetry. Although very heavy $B-L$ gauge boson $M_{Z_{B-L}} \gtrsim 10^3 M_{N_1}$ is required for suppress scattering effects in many cases, $M_{Z_{B-L}} \sim 3 M_{N_1}$ is possible in our model because CP asymmetry is large, $\epsilon_1 \sim 0.1$.

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