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Gauged $B - 3L_\tau$ and Baryogenesis

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Abstract

It has recently been shown that by extending the minimal standard model to include a right-handed partner to ν_τ, it is possible to gauge the $B - 3L_\tau$ quantum number consistently. If we add two scalar triplets, one trivial (ξ_1) and one nontrivial (ξ_2) under $B - 3L_\tau$, it is possible also to have desirable neutrino masses and mixing for neutrino oscillations. At the same time, a lepton asymmetry can be generated in the early universe through the novel mechanism of the decay of the heavier ξ_1 into the lighter ξ_2 plus a neutral singlet (ζ^0). This lepton asymmetry then gets converted into a baryon asymmetry at the electroweak phase transition.

It has recently been shown[1] that the minimal standard $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge model of quarks and leptons may be extended to include an anomaly-free gauge factor $U(1)_{B-3L_\tau}$ if ν_τ has a right-handed singlet partner, but not ν_e or ν_μ. The scale of symmetry breaking of this $B-3L_\tau$ gauge group may even be lower[2] than that of electroweak symmetry breaking. In the minimal standard model, there are dimension-six baryon-number nonconserving operators[3] of the form Q^3L which would induce the proton to decay. In the $B-3L_\tau$ gauge model, the lowest dimensional baryon-number nonconserving operator is of the form Q^9L_τ, hence such processes are suppressed by 22 powers of some higher energy scale and become totally negligible. To obtain a baryon asymmetry of the universe in this model, it is natural to propose instead that a primordial lepton asymmetry is generated[4, 5] which then gets converted into a baryon asymmetry during the electroweak phase transition[6].

In order to generate a lepton asymmetry in the early universe, we must have lepton-number nonconserving interactions which also violate C and CP, as well as the existence of an epoch when such processes are out of thermal equilibrium[7]. The canonical way[4] of achieving this is to use heavy right-handed singlet neutrinos which also allow the known left-handed neutrinos $(\nu_e, \nu_\mu, \nu_\tau)$ to acquire small seesaw masses. An equally attractive scenario was recently proposed[5] where heavy Higgs triplets are used. In the $B-3L_\tau$ gauge model, since ν_e and ν_μ have no right-handed singlet partners, the natural thing to do is to adopt a variation of the latter mechanism. In fact, as we see below, the requirement of a desirable neutrino mass matrix and the absence of an unwanted pseudo-Goldstone boson together imply a successful leptogenesis scenario in this model without any further extension.

In the standard model, a general neutrino mass matrix may be obtained with the addition of one heavy Higgs triplet, but to generate a lepton asymmetry, two such triplets are required[5]. In the minimal $B-3L_\tau$ gauge model, only ν_τ gets a mass. Furthermore, because $B-3L_\tau$ is a gauge symmetry, it is not obvious *a priori* how that will affect the conversion of a primordial $L_e + L_\mu$ asymmetry into a baryon asymmetry during the electroweak phase transition. In this paper we will address both issues, *i.e.* neutrino masses and baryogenesis, and show how they may have a common solution.

The fermion content of our model is identical to that of the original model[1]. The quarks and leptons transform under $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-3L_\tau}$ as follows:

$$\begin{pmatrix} u_i \\ d_i \end{pmatrix}_L \sim (3,2,1/6;1/3), \quad u_{iR} \sim (3,1,2/3;1/3), \quad d_{iR} \sim (3,1,-1/3;1/3); \tag{1}$$

$$\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L, \begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L \sim (1,2,-1/2;0), \quad e_R, \mu_R \sim (1,1,-1;0); \tag{2}$$

$$\begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L \sim (1,2,-1/2;-3), \quad \tau_R \sim (1,1,-1;-3), \quad \nu_{\tau R} \sim (1,1,0;-3). \tag{3}$$

The $U(1)_{B-3L_\tau}$ gauge boson X does not couple to e or μ or their corresponding neutrinos. It can thus escape detection in most experiments. Although it does couple to quarks as in a previously proposed model[8], such signatures are normally overwhelmed by the enormous quantum-chromodynamics (QCD) background. On the other hand, X does couple to τ and ν_τ, so it could be observed through its decay into $\tau^+\tau^-$ pairs[2]. Furthermore, the ν_τ-quark interactions in this scenario may affect the oscillations of neutrinos inside the sun and the earth and contribute[9] to the zenith-angle dependence of the atmospheric neutrino deficit[10]. Since the interactions of X violates $e-\mu-\tau$ universality, present experimental constraints limit its coupling and mass[2]. For example, $g_X < 0.1$ is required for $m_X < 50$ GeV.

The minimal scalar sector of this model consists of the standard Higgs doublet,

$$\begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix} \sim (1,2,1/2;0) \tag{4}$$

which breaks the electroweak gauge symmetry $SU(2)_L \times U(1)_Y$ down to $U(1)_{em}$ and a neutral singlet,

$$\chi^0 \sim (1,1,0;6) \tag{5}$$

which couples to $\nu_{\tau R}\nu_{\tau R}$ and breaks the $U(1)_{B-3L_\tau}$ gauge symmetry. The resulting theory allows $\nu_{\tau L}$ to obtain a seesaw mass[11] of order 1 eV and retains B as an additively conserved quantum number and L_τ as a multiplicatively conserved quantum number.

In the original model[1], the other two neutrinos (ν_e, ν_μ) acquire masses and mix with ν_τ radiatively. In our present work, we propose instead to use the mechanism of Ref.[5] and add a couple of Higgs triplets:

$$\begin{pmatrix} \xi_1^{++} \\ \xi_1^{+} \\ \xi_1^{0} \end{pmatrix} \sim (1,3,1;0) \quad \text{and} \quad \begin{pmatrix} \xi_2^{++} \\ \xi_2^{+} \\ \xi_2^{0} \end{pmatrix} \sim (1,3,1;3). \tag{6}$$

Now ξ_1 will give small masses to ν_e and ν_μ, and will also generate a ξ_2-asymmetry of the universe when it decays at a very high temperature. Furthermore, ξ_2 will mix ν_e and ν_μ with ν_τ, and its interactions will convert the ξ_2-asymmetry into a $L_e + L_\mu$ asymmetry of the universe. Since $B - 3L_\tau$ may well be an unbroken gauge symmetry during the electroweak phase transition, there may not be any $B - 3L_\tau$ asymmetry. In that case, the total $B - L$ asymmetry is the same as the $-(L_e + L_\mu)$ asymmetry, which will get converted into the baryon asymmetry of the universe during the electroweak phase transition.

The above scalar sector contains a pseudo-Goldstone boson which comes about because there are 3 global $U(1)$ symmetries in the Higgs potential and only 2 local $U(1)$ symmetries which get broken. In addition, there is no CP-violating complex phase which is necessary for generating a lepton asymmetry of the universe. However, if an extra neutral scalar $\zeta^0 \sim (1,1,0,-3)$ is added, then the Higgs potential will have additional terms which eliminate the unwanted global $U(1)$ symmetry, and one of these couplings will also have to be complex, allowing thus enough CP violation to generate a lepton asymmetry, as explained below.

The triplet scalar fields $\xi_a,\ (a = 1,2)$ do not acquire any vacuum expectation value (vev)

to start with. At the tree level, we can write down the relevant part of the Lagrangian as

$$
\begin{aligned}
-\mathcal{L} \;=\; & \sum_{a=1,2} M_a^2 \xi_a^\dagger \xi_a + \sum_{i,j=e,\mu} f_{1ij}[\xi_1^0 \nu_i \nu_j + \xi_1^+(\nu_i l_j + l_i \nu_j)/\sqrt{2} + \xi_1^{++} l_i l_j] \\
& + \sum_{i=e,\mu} f_{2i\tau}[\xi_2^0 \nu_i \nu_\tau + \xi_2^+(\nu_i \tau + l_i \nu_\tau)/\sqrt{2} + \xi_2^{++} l_i \tau] + f_\tau \nu_{\tau R} \nu_{\tau R} \chi^0 \\
& + \mu_1[\bar{\xi}_1^0 \phi^0 \phi^0 + \sqrt{2}\xi_1^- \phi^+ \phi^0 + \xi_1^{--} \phi^+ \phi^+] + \mu_2 \xi_1^\dagger \xi_2 \zeta^0 + g \zeta^0 \zeta^0 \chi^0 \\
& + h \zeta^{0*}[\bar{\xi}_2^0 \phi^0 \phi^0 + \sqrt{2}\xi_2^- \phi^+ \phi^0 + \xi_2^{--} \phi^+ \phi^+] + h.c. \qquad (7)
\end{aligned}
$$

Interactions of the scalar triplets $\xi_a,\ (a = 1, 2)$ break the L_e and L_μ numbers explicitly, whereas L_τ is conserved. Since there is no spontaneous breaking of the lepton numbers, there are no massless Goldstone bosons (majorons) which can contribute to the invisible width of the Z boson. After electroweak symmetry breaking when the doublet acquires a nonzero *vev*, and the breaking of $B-3L_\tau$, there will be induced *vev*'s for these scalar triplets,

$$
\langle \xi_1^0 \rangle \simeq \frac{-\mu_1 v^2}{M_1^2} \qquad \text{and} \qquad \langle \xi_2^0 \rangle \simeq \frac{-h \langle \zeta^0 \rangle v^2}{M_2^2}.
$$

Since the masses of these scalar triplets and the would-be majorons are very large, they cannot contribute to the width of the Z boson.

At low energies, we can integrate out the heavier triplet fields and write down an effective neutrino mass matrix in the basis $\{\bar{\nu}_{eL} \;\; \bar{\nu}_{\mu L} \;\; \bar{\nu}_{\tau L} \;\; \nu_{\tau R}\}$ as,

$$
M_\nu = \begin{pmatrix} f_{1ee}\langle \xi_1^0 \rangle & f_{1e\mu}\langle \xi_1^0 \rangle & f_{2e\tau}\langle \xi_2^0 \rangle & 0 \\ f_{1e\mu}\langle \xi_1^0 \rangle & f_{1\mu\mu}\langle \xi_1^0 \rangle & f_{2\mu\tau}\langle \xi_2^0 \rangle & 0 \\ f_{2e\tau}\langle \xi_2^0 \rangle & f_{2\mu\tau}\langle \xi_2^0 \rangle & 0 & m_\tau^D \\ 0 & 0 & m_\tau^D & f_\tau \langle \chi^0 \rangle \end{pmatrix} \qquad (8)
$$

where m_τ^D is the Dirac mass term for ν_τ and $f_\tau\langle \chi^0 \rangle$ is the Majorana mass of $\nu_{\tau R}$. The left-handed $\nu_{\tau L}$ will then get a seesaw mass, which can be of order 1 eV. The out-of-equilibrium condition for the generation of a lepton asymmetry dictates that the mass of the triplet ξ_1 be of order 10^{13} GeV, which implies that the e and μ mass elements are also of order 1 eV or less. As we will see later, ν_e and ν_μ may mix with ν_τ to form a desirable phenomenological mass matrix for neutrino oscillations. The constraints on these elements come from the consideration of a realistic baryon asymmetry of the universe as we show below.

Around the time of the electroweak phase transition, we assume that $B-3L_\tau$ is conserved; hence there cannot be any $B-3L_\tau$ asymmetry. In particular, L_τ is exactly conserved. However, L_e and L_μ numbers are broken explicitly at some high scale M_1 in the decays of the triplets ξ_1. At such high energies, $SU(2)_L$ gauge invariance means that we need only consider one of its components, say ξ_1^{++}, which has the following decay modes:

$$\xi_1^{++} \rightarrow \begin{cases} l_i^+ l_j^+ & (L_e + L_\mu = -2;\ n_{\xi_2} = 0) \\ \xi_2^{++}\zeta^0 & (L_e + L_\mu = -1;\ n_{\xi_2} = 1) \\ \phi^+\phi^+ & (L_e + L_\mu = 0;\ n_{\xi_2} = 0). \end{cases} \tag{9}$$

Here we assumed that most of the time $\xi_2^{++} \rightarrow l_i^+ l_\tau^+$, and the other decay mode of ξ_2 never comes to equilibrium so that $L_e + L_\mu = -1$ for ξ_2. We will discuss this point later. In the following we will first explain how ξ_1 decay generates a ξ_2 asymmetry.

The first decay mode does not play any role in the generation of a lepton asymmetry. Here CP violation comes from the interference of the tree-level and one-loop diagrams of Figures 1 and 2. The scalar potential has one CP-violating phase in the product $\mu_1^*\mu_2 h$, which cannot be absorbed by redefinitions and produces a ξ_2-asymmetry when ξ_1 decays. As a result, the decays of ξ_1^{++} and ξ_1^{--} will create more ξ_2^{++} than ξ_2^{--} or vice versa; hence a ξ_2-asymmetry $\delta = (n_{\xi_2} - n_{\xi_2^\dagger})/n_\gamma$ will be created, given by

$$\delta \simeq \frac{Im\,[\mu_1^*\mu_2 h]}{16\pi^2 g_* M_1^2}\left[\frac{M_1}{\Gamma_1}\right], \tag{10}$$

where g_* is the total number of relativistic degrees of freedom and

$$\Gamma_1 = \frac{1}{8\pi}\left(\frac{|\mu_1|^2 + |\mu_2|^2}{M_1} + \sum_{i,j}|f_{1ij}|^2 M_1\right) \tag{11}$$

is the decay rate of the triplet ξ_1. This ξ_2-asymmetry will also have an apparent charge asymmetry, which will be compensated by an asymmetry in ϕ^+ and ϕ^-. In earlier models of leptogenesis, a lepton asymmetry is generated when the heavy particles decay into light leptons and CP violation enters in the vertex corrections[4] or in the mass matrix[5, 12]. In contrast, we generate in the present scenario an asymmetry in ξ_2 through the quartic scalar couplings, which then generate a lepton asymmetry.

For the generation of the ξ_2-asymmetry, this decay rate should also satisfy the out-of-equilibrium condition[13]

$$\Gamma_1 < \sqrt{1.7 g_*} \frac{T^2}{M_{Pl}} \qquad \text{at} \;\; T = M_1, \tag{12}$$

where M_{Pl} is the Planck scale. We assume $M_1 \gg M_2$, so that when M_1 decays, ξ_2 is essentially massless. Taking $\mu_{1,2}/M_1 \sim 0.1$, $f_{1ij} \sim 1$, the out-of-equilibrium condition is satisfied with $M_1 > 10^{14}$ GeV. However, even if we choose $M_1 \sim 10^{13}$ GeV, the generated ξ_2-asymmetry will be only less by a factor $S \sim 10^{-2}$, which is still large enough to explain the baryon asymmetry of the universe for a value of $h \sim 10^{-4}$. This gives us the e and μ neutrino mass matrix elements to be of order 1 eV.

At a temperature $T < M_1$, there will be a ξ_2-asymmetry. The decays of ξ_2 also break lepton number,

$$\xi_2^{++} \rightarrow \begin{cases} l_i^+ l_\tau^+ & (L_e + L_\mu = -1) \\ \phi^+ \phi^+ \zeta^{0*} & (L_e + L_\mu = 0). \end{cases} \tag{13}$$

If both of these decay modes are in equilibrium at any time, that will erase the lepton asymmetry of the universe[6, 13, 14]. We must therefore require that at least one of these interactions and the scattering process,

$$l_i^+ l_\tau^+ \rightarrow \phi^+ \phi^+ \zeta^{0*}$$

to satisfy the out-of-equilibrium condition till the electroweak symmetry breaking phase transition is over.

For the choice $h \sim 10^{-4}$, we may take $M_2 > 10^5$ GeV to ensure that

$$\Gamma_2(\xi_2 \rightarrow \phi\phi\zeta^*) = \frac{h^2}{16\pi^2} \frac{M_2}{8\pi} < \sqrt{1.7 g_*} \frac{T^2}{M_{Pl}} \qquad \text{at } T \geq M_2 \tag{14}$$

so that ξ_2 can hardly decay into three scalars at any time. However, we would like the other decay mode of ξ_2 to be fast, so that the ξ_2-asymmetry generated during the ξ_1 decay gets converted into a lepton asymmetry. In other words, since the number of ξ_2 is different from the numbers of $\xi_2^\dagger$, the number of leptons generated in decays of ξ_2 will be different from the number of antileptons generated in decays of $\xi_2^\dagger$.

We take $f_{2i} \sim 0.1$, so that the two-lepton decay mode of ξ_2 is in equilibrium for most of the time,

$$\Gamma_2(\xi_2 \to l_i l_\tau) = \frac{\sum_i |f_{2i\tau}|^2}{8\pi} M_2 > \sqrt{1.7 g_*} \frac{T^2}{M_{Pl}} \qquad \text{during } M_c \geq T \geq M_2 \tag{15}$$

where $M_c \simeq 10^9$ GeV. During this period from M_c to M_2, the ξ_2-asymmetry will get converted into a $L_e + L_\mu$ asymmetry of the universe. The interaction $\xi_2 L_i L_\tau$ will also be in equilibrium, which will relate their chemical potentials: $\mu_{\xi_2} = \mu_{L_i} + \mu_{L_\tau}$ (notations will be explained later). The generated $L_e + L_\mu$ asymmetry is accompanied by an equal amount of L_τ asymmetry. However, that is compensated exactly by the ζ-asymmetry created at the time of ξ_1 decay. This ζ-asymmetry generates an equal and opposite amount of L_τ asymmetry through $\zeta + \zeta \to \chi^* \to \nu_{\tau R} + \nu_{\tau R}$, which compensates the L_τ asymmetry in ξ_2 decay. As a result, there will not be any net L_τ asymmetry as expected, since $B - 3L_\tau$ is exactly conserved at this time. The generated L_e and L_μ asymmetries together give the $B - L$ asymmetry

$$n_L = \frac{1}{2} \delta S.$$

The above choice of parameter values will give us the neutrino mass mixing of the e and the μ to the τ neutrinos of order 1 eV.

We will now see how this lepton asymmetry can get converted into the baryon asymmetry of the universe during the electroweak phase transition[6]. We consider all the particles to be ultrarelativistic, which is the case above the electroweak scale, but at lower energies, although we understand that a careful analysis has to include the mass corrections, we ignore them since they are small and cannot change the conclusion drastically. The particle asymmetry, *i.e.* the difference between the number of particles (n_+) and the number of antiparticles (n_-) can be given in terms of the chemical potential of the particle species μ (for antiparticles the chemical potential is $-\mu$) as

$$n_+ - n_- = n_d \frac{gT^3}{6} \left(\frac{\mu}{T} \right), \tag{16}$$

where $n_d = 2$ for bosons and $n_d = 1$ for fermions.

In the rest of this discussion we will assume that after the triplets ξ_1 and ξ_2 have decayed, enough lepton asymmetry was generated. This will give nonvanishing $\mu_{\nu e}$ and $\mu_{\nu\mu}$, which are directly related to n_L. When these neutrinos interact with other particles in equilibrium, the chemical potentials get related by simple additive relations, and that will allow us to relate this lepton asymmetry n_L to the baryon asymmetry during the electroweak phase transition.

At energies near the electroweak phase transition, most of the interactions are in equilibrium. These include the sphaleron[15] induced electroweak $B + L$ violating interaction arising due to the nonperturbative axial-vector anomaly[16]. In Table 1, we give the interactions and the corresponding relations between the chemical potentials. In the third column we give the chemical potential which we eliminate using the given relation. We start with chemical potentials of all the quarks ($\mu_{uL}, \mu_{dL}, \mu_{uR}, \mu_{dR}$); the e and μ leptons ($\mu_{aL}, \mu_{\nu aL}, \mu_{aR}$, where $a = e, \mu$); the τ leptons ($\mu_{\tau L}, \mu_{\nu\tau L}, \mu_{\tau R}, \mu_{\nu\tau R}$); the gauge bosons ($\mu_W$ for W^-, and 0 for all others); and the Higgs scalars ($\mu_-^\phi, \mu_0^\phi, \mu^\chi, \mu^\zeta$). The triplets have decayed away much before the electroweak phase transition and have decoupled; hence they do not contribute to the present analysis.

We can then express all the chemical potentials in terms of the following independent chemical potentials only,

$$\mu_0 = \mu_0^\phi; \quad \mu_W; \quad \mu_u = \mu_{uL}; \quad \mu_a = \mu_{\nu eL} = \mu_{\nu\mu L}; \quad \mu_\tau = \mu_{\nu\tau L}. \tag{17}$$

Table 1: Relations among the chemical potentials

Interactions	μ relations	μ eliminated
$D_\mu \phi^\dagger D_\mu \phi$	$\mu_W = \mu_-^\phi + \mu_0^\phi$	μ_-^ϕ
$\overline{q_L} \gamma_\mu q_L W^\mu$	$\mu_{dL} = \mu_{uL} + \mu_W$	μ_{dL}
$\overline{l_L} \gamma_\mu l_L W^\mu$	$\mu_{iL} = \mu_{\nu iL} + \mu_W$	$\mu_{iL}, i = e, \mu, \tau$
$\overline{q_L} u_R \phi^\dagger$	$\mu_{uR} = \mu_0 + \mu_{uL}$	μ_{uR}
$\overline{q_L} d_R \phi$	$\mu_{dR} = -\mu_0 + \mu_{dL}$	μ_{dR}
$\overline{l_{aL}} e_{aR} \phi$	$\mu_{aR} = -\mu_0 + \mu_{aL}$	$\mu_{aR}, a = e, \mu$
$\overline{l_{\tau L}} e_{\tau R} \phi$	$\mu_{\tau R} = -\mu_0 + \mu_{\tau L}$	$\mu_{\tau R}$
$\overline{\nu_{\tau R}}^c \nu_{\tau R} \chi$	$\mu^\chi = -2\mu_{\nu\tau R}$	μ^χ
$\overline{l_{\tau L}} \nu_{\tau R} \phi^\dagger$	$\mu_{\nu\tau R} = \mu_0 + \mu_{\tau L}$	$\mu_{\nu\tau R}$
$\zeta\zeta\chi^0$	$\mu^\chi = 2\mu^\zeta$	μ^ζ

We can further eliminate one of these five potentials by making use of the relation given by the sphaleron processes. Since the sphaleron interactions are in equilibrium, we can write down the following $B + L$ violating relation among the chemical potentials for three generations,

$$9\mu_u + 6\mu_W + 2\mu_a + \mu_\tau = 0. \tag{18}$$

We then express the baryon number, lepton numbers and the electric charge and the hypercharge number densities in terms of these independent chemical potentials,

$$B = 12\mu_u + 6\mu_W \tag{19}$$

$$L_e = L_\mu = 3\mu_a + 2\mu_W - \mu_0 \tag{20}$$

$$L_\tau = 4\mu_\tau + 2\mu_W \tag{21}$$

$$Q = 24\mu_u + (12 + 2m)\mu_0 - (4 + 2m)\mu_W \tag{22}$$

$$Q_3 = -(10 + m)\mu_W \tag{23}$$

where m is the number of Higgs doublets ϕ.

At temperatures above the electroweak phase transition, $T > T_c$, both Q and Q_3 must vanish. In addition, since $B - 3L_\tau$ is also a gauge symmetry, this charge must also vanish. These three conditions and the sphaleron induced $B-L$ conserving, $B+L$ violating condition can be expressed as

$$
\begin{aligned}
< Q >= 0 \quad &\Longrightarrow \quad \mu_0 = \frac{-12}{6+m}\mu_u && (24) \\
< Q_3 >= 0 \quad &\Longrightarrow \quad \mu_W = 0 && (25) \\
< B - 3L_\tau >= 0 \quad &\Longrightarrow \quad \mu_\tau = \mu_u && (26) \\
\text{Sphaleron transition} \quad &\Longrightarrow \quad \mu_a = -5\mu_u && (27)
\end{aligned}
$$

Using these relations we can now write down the baryon number, lepton number, and their combinations in terms of the $B - L$ number density, which remains invariant under all electroweak phase transitions. They are

$$
\begin{aligned}
B &= \frac{36+6m}{102+19m}(B-L) && (28) \\
L_e = L_\mu &= \frac{-78-15m}{204+38m}(B-L) && (29) \\
L_\tau &= \frac{12+2m}{102+19m}(B-L) && (30) \\
B+L &= \frac{-30-7m}{102+19m}(B-L) && (31)
\end{aligned}
$$

We will now consider two possibilities. In the first, the $B-3L_\tau$ gauge symmetry is broken after the electroweak phase transition. Then, at temperatures below the electroweak phase transition, the other relations remain the same, but it is no longer neccessary to make Q_3 vanishing. Since ϕ acquires a *vev*, we require $\mu_0 = 0$. With this change we can now relate all the chemical potentials in terms of μ_u as

$$
\begin{aligned}
< Q >= 0 \quad &\Longrightarrow \quad \mu_W = \frac{12}{2+m}\mu_u && (32) \\
\langle \phi_0 \rangle \neq 0 \quad &\Longrightarrow \quad \mu_0 = 0 && (33)
\end{aligned}
$$

$$< B - 3L_\tau >= 0 \implies \mu_\tau = \mu_u \tag{34}$$

$$\text{Sphaleron transition} \implies \mu_a = -3\mu_W - 5\mu_u \tag{35}$$

This will then allow us to write down the baryon number, lepton number, and their combinations in terms of the $B - L$ number density as

$$B = \frac{48 + 6m}{146 + 19m}(B - L) \tag{36}$$

$$L_e = L_\mu = \frac{-114 - 15m}{292 + 38m}(B - L) \tag{37}$$

$$L_\tau = \frac{16 + 2m}{146 + 19m}(B - L) \tag{38}$$

$$B + L = \frac{-50 - 7m}{292 + 38m}(B - L) \tag{39}$$

Thus after the electroweak phase transition, the $L_e + L_\mu$ asymmetry n_L generated after the scalar triplets ξ_1 and then ξ_2 have decayed, which is equal to the $B - L$ asymmetry at that time, will get converted into a B asymmetry during the electroweak phase transition. Although any existing $B + L$ asymmetry gets washed out, we still get a non-zero $B + L$ asymmetry after the electroweak phase transition from the same $B - L$ asymmetry. For consistency we check that the $B - L$ asymmetry remains the same during the electroweak phase transition and there is no $B - 3L_\tau$ asymmetry.

We will now consider the other possibility when $B - 3L_\tau$ is broken before the electroweak phase transition. In this case the electroweak symmetry is unbroken and we still have $< Q_3 >= 0$, but $< B - 3L_\tau >\neq 0$ so that $\mu_\chi = 0$. The constraints in this case are

$$< Q >= 0 \implies \mu_0 = \frac{-12}{6 + m}\mu_u \tag{40}$$

$$< Q_3 >= 0 \implies \mu_W = 0 \tag{41}$$

$$\langle \chi \rangle \neq 0 \implies \mu_\chi = 0 \Rightarrow \mu_{\nu_\tau R} = 0;\ \mu_\tau = -\mu_0 \tag{42}$$

$$\text{Sphaleron transition} \implies \mu_a = \mu_0 - \frac{9}{2}\mu_u \tag{43}$$

The baryon and lepton asymmeties are now related by

$$
\begin{aligned}
B &= \frac{24+4m}{66+13m}(B-L) && (44)\\
L_e = L_\mu &= \frac{-58-9m}{132+26m}(B-L) && (45)\\
L_\tau &= \frac{16}{66+13m}(B-L) && (46)\\
B+L &= \frac{-18-5m}{66+13m}(B-L) && (47)
\end{aligned}
$$

Finally we give the relations between the baryon and lepton asymmetries after both the electroweak symmetry and the $B-3L_\tau$ symmetry are broken. The final asymmetry will not depend on whether the electroweak symmetry was broken before or after the $B-3L_\tau$ symmetry is broken. Now we have $\mu_\chi = 0$ but $< B-3L_\tau > \neq 0$ and $\langle\phi\rangle = 0 \Rightarrow \mu_0 = 0 \Rightarrow \mu_\tau = 0$. We can then express the other chemical potentials in terms of μ_u as

$$
\begin{aligned}
< Q >= 0 &\implies \mu_W = \frac{12}{2+m}\mu_u && (48)\\
\langle\phi\rangle \neq 0 &\implies \mu_0 = 0 && (49)\\
\langle\chi\rangle \neq 0 &\implies \mu_\chi = 0 \Rightarrow \mu_{\nu_\tau R} = 0;\ \mu_\tau = -\mu_0 = 0 && (50)\\
\text{Sphaleron transition} &\implies \mu_a = -3\mu_W - \frac{9}{2}\mu_u && (51)
\end{aligned}
$$

which then let us write the baryon and lepton numbers as some combinations of $B-L$ as

$$
\begin{aligned}
B &= \frac{32+4m}{98+13m}(B-L) && (52)\\
L_e = L_\mu &= \frac{-74-9m}{196+26m}(B-L) && (53)\\
L_\tau &= \frac{8}{98+13m}(B-L) && (54)\\
B+L &= \frac{-34-5m}{98+13m}(B-L) && (55)
\end{aligned}
$$

The final baryon asymemtry of the universe is about 1/3 that of the $B-L$ asymmetry. Hence in the present scenario, the generated $B-L$ asymmetry n_L will get converted into a baryon asymmetry after the electroweak and $B - 3L_\tau$ symmetries are broken, and the present baryon asymmetry of the universe will be given by

$$n_b \sim \frac{1}{6} S\delta \tag{56}$$

which is of the order of 10^{-10}, for the choice of parameter values we have considered earlier.

To summarize, we studied an extension of the standard model to include a $B - 3L_\tau$ gauge symmetry, which may be broken below the electroweak symmetry breaking scale. The Higgs structure is modified to explain the baryon asymmetry of the universe, which comes about in an unconventional way. We first generate an asymmetry in the number of scalars (n_{ξ_2}), through only scalar interactions. This n_{ξ_2} generates the $L_e + L_\mu$ asymmetry when these scalars decay. During the electroweak phase transition, the latter gets converted into a baryon asymmetry of the universe. Neutrino masses and mixing are obtained naturally in this scenario.

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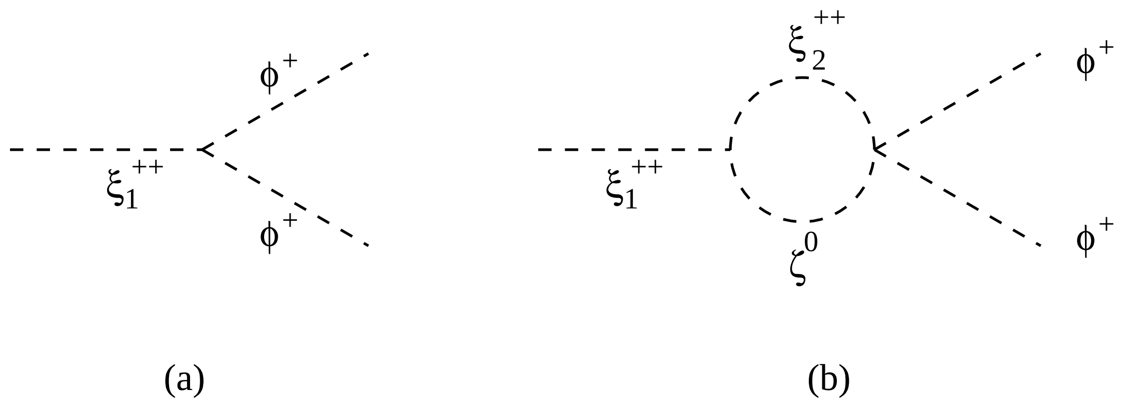

Figure 1: Tree-level and one-loop diagrams for $\xi_1 \to \phi\phi$.

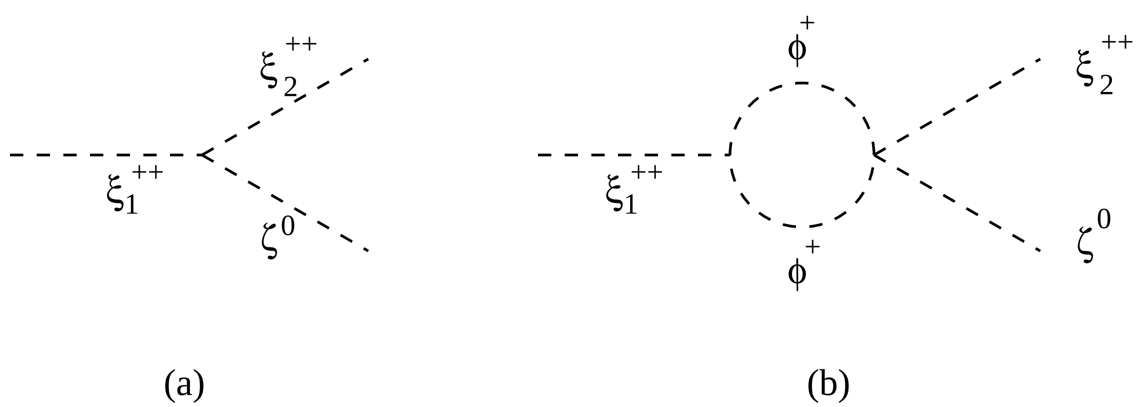

Figure 2: Tree-level and one loop diagrams for $\xi_1 \to \xi_2\zeta$.