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Cogenesis in a universe with vanishing $B-L$ within a gauged $U(1)_x$ extension

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ABSTRACT

We consider a gauged $U(1)_x$ extension of the standard model and of the minimal supersymmetric standard model where the dark matter fields are charged under $U(1)_x$ and carry lepton number while the standard model fields and fields of the minimal supersymmetric standard model are neutral under $U(1)_x$. We consider leptogenesis in this class of models with all fundamental interactions having no violation of lepton number, and the total $B-L$ in the universe vanishes. Such leptogenesis leads to equal and opposite lepton numbers in the visible sector and in the dark sector, and thus also produces asymmetric dark matter. Part of the lepton number generated in the leptonic sector subsequently transfer to the baryonic sector via sphaleron interactions. The stability of the dark particles is protected by the $U(1)_x$ gauge symmetry. A kinetic mixing between the $U(1)_x$ and the $U(1)_Y$ gauge bosons allows for dissipation of the symmetric component of dark matter. The case when $U(1)_x$ is $U(1)_{B-L}$ is also discussed for the supersymmetric case. This case is particularly interesting in that we have a gauged $U(1)_{B-L}$ which ensures the conservation of $B-L$ with an initial condition of a vanishing $B-L$ in the universe. Phenomenological implications of the proposed extensions are discussed, which includes implications for electroweak physics, neutrino masses and mixings, and lepton flavor changing processes such as $\ell_i \to \ell_j \gamma$. We also briefly discuss the direct detection of the dark matter in the model.

1. Introduction

Three of the important puzzles in cosmology relate to the origin of baryon asymmetry in the Universe, the nature of dark matter and the cosmic coincidence. Thus the visible universe exhibits an excess of baryons over anti-baryons and this excess is often displayed as the baryon number density to the entropy density ratio [1]

$$B/s \sim 6 \times 10^{-10}. \tag{1}$$

The basic tenets of how to generate baryon (lepton) excess has been known since the work of Sakharov [2], and consist of three conditions, i.e., the existence of baryon (or lepton) number violation, the presence of C and CP violating interactions, and out of equilibrium processes. In the standard model the ratio B/s is computed to be too small to fit observation pointing to the existence of beyond the standard model physics. Standard model also does not provide us with a candidate for dark matter and the astrophysical evidence for its presence again points to the existence of new physics beyond the standard model. Additionally one has the cosmic coincidence puzzle, i.e., the fact that the amount of dark matter and the amount of visible matter in the Universe are comparable. Specifically one has [3]

$$\frac{\Omega_{\rm DM} h_0^2}{\Omega_{\rm B} h_0^2} \approx 5.5. \tag{2}$$

The comparable sizes of the amounts of dark matter and of visible matter point to the possibility of a common origin of the two. This can be explained by the so-called asymmetric dark matter hypothesis where the dark particles are in thermal equilibrium with the standard model particles or with the particles of the minimal supersymmetric standard model in the early universe, and thus their chemical potentials are of the same order. The satisfaction of Eq. (2) then occurs via a constraint on the dark matter mass (for a sample of recent works see [4,5] and for reviews see [6]). Alternative schemes where dark matter carrying a lepton number (or a baryon number) is created first and a portion of it subsequently transfer to the visible sector have been considered in [7,8].

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Cogenesis of baryon/lepton asymmetry and the asymmetric dark matter have also been discussed recently in [9].

An important constraint on model building is the requirement that dark matter be stable, i.e., the dark particles does not decay into lighter standard model particles. In this work we consider an extension of the standard model and of the minimal supersymmetric standard model where the dark fields are charged under a $U(1)_x$ gauge symmetry while the standard model fields are neutral under $U(1)_x$, which forbids dark particles decay into the standard model particles and thus guarantees the stability of the dark matter. Additionally, the asymmetry of the dark particles generated in the early universe will not be washed out by Majorana mass terms since they are forbidden by the $U(1)_x$ gauge symmetry.[1] In the supersymmetric case, a gauged $U(1)_{B-L}$ model is also discussed.

Most conventional models of baryogenesis or leptogenesis assume that the fundamental vertices violate either baryon number or lepton number or both in conformity with the first Sakharov condition [11,12]. However, in this work we consider leptogenesis where the fundamental interactions conserve lepton number and leptogenesis consists in generating equal and opposite lepton numbers in the visible and in the dark sectors. Subsequently the sphaleron processes transmute a part of the leptons into baryons. The total $B–L$ in the universe is exactly conserved. This mechanism bypasses the difficulty in the GUT baryogenesis where a vanishing total $B–L$ implies that the baryon asymmetry generated would be washed out by the sphaleron interactions. While this idea has been recently pursued by several authors [13,14], our analysis differs significantly in structure and in content from previous works [13,14].[2] A more detailed comparison with these works is given at the end of Section 4. Earlier works on Dirac leptogenesis [16] can also generate the asymmetry in the visible sector starting from a $B–L$ vanishing universe. Due to the tiny Yukawa coupling, right-handed (Dirac) neutrinos would not be in thermal equilibrium with left-handed neutrinos, hence the sphaleron interactions which operate only on $SU(2)$ fields, will not wash out the lepton number stored in the right-handed neutrinos and thus the asymmetry is created.

The outline of the rest of the Letter is as follows: In Section 2 we discuss leptogenesis and the generation of asymmetric dark matter in a non-supersymmetric model where the vertices have no lepton number violation. The dark matter consists of two fermionic fields which carry the same lepton numbers but opposite $U(1)_x$ charges. Here we also compute the mass of the dark particles which satisfy the cosmic coincidence of Eq. (2). In Section 3 we extend the analysis to the supersymmetric case. The main difference in the analysis of Section 3 from the analysis of Section 2 is that in the supersymmetric case there are more species of dark matter particles. Specifically we have four types of fermionic particles and their bosonic super-partners which carry different combinations of the lepton numbers and $U(1)_x$ charges. We also discuss the possibility that $U(1)_x$ is $U(1)_{B-L}$. In Section 4 we discuss the phenomenology related to these models. Conclusions are given in Section 5.

[1] Models which allow a Majorana term for dark matter can undergo oscillations where the dark particle oscillates to its anti-particle. Such processes over the lifetime of the universe can produce symmetric dark matter which can lead to pair annihilation and wipe out the asymmetric dark matter [10].

[2] Baryogenesis with a gauged $U(1)_B$ symmetry is discussed in [15], where the dark sector and the visible sector carry the opposite baryon number and the total baryon number in the universe is conserved. While in this work a pre-existing excess of lepton number has been assumed, thus the total $B–L$ in the universe is not vanishing.

2. Non-supersymmetric model

We begin by considering the set of fields N_i, ψ, ϕ, X, X' with lepton number assignments $(0, +1, -1, +1/2, +1/2)$. Here N_i ($i \geqslant 2$) are Majorana fermions, ψ, X, X' are Dirac fields and ϕ is a complex scalar field. The fields N_i, ψ, ϕ are heavy and will decay into lighter fields and eventually disappear and there would be no vestige left of them in the current universe. The dark sector is constituted of two fermionic fields X, X', which as indicated above each carry a lepton number $+1/2$ and are oppositely charged under the dark sector gauge group $U(1)_x$ with gauge charges $(+1, -1)$. All other fields are neutral under $U(1)_x$. We assume their interactions to have the following form which conserve both the lepton number and the $U(1)_x$ gauge symmetry:

$$\mathcal{L} = \lambda_i \bar{N}_i \psi \phi + \beta \bar{\psi} L H + \gamma \phi \bar{X}^c X' + h.c., \tag{3}$$

where the couplings λ_i are assumed to be complex and the couplings β, γ are assumed to be real. In addition we add mass terms so that

$$-\mathcal{L}_m = M_i \bar{N}_i N_i + m_1 \bar{\psi} \psi + m_2^2 \phi^* \phi + m_X \bar{X} X + m_{X'} \bar{X}' X'. \tag{4}$$

Here N_i have Majorana masses, while ψ, X, X' have Dirac masses. We assume the mass hierarchy $M_i \gg m_1 + m_2$, $m_1 \sim m_2 \gg m_X + m_{X'}$. We will see later that $m_X, m_{X'}$ are around 1 GeV. Consistent with the above constraint, m_1, m_2 which are the masses of ψ and ϕ respectively, could span a wide range from order of TeV to much higher scales.

In the early universe, the out-of-equilibrium decays of the heavy Majorana fields N_i produce a heavy Dirac field ψ and a heavy complex scalar field ϕ. The CP violation due to the complex couplings λ_i generates an excess of ψ, ϕ over their anti-particles $\bar{\psi}, \phi^*$ which carry the opposite lepton numbers. Since the lepton number carried by ψ and ϕ always sums up to zero, the out-of-equilibrium decays of N_i do not generate an excess of lepton number in the universe. Further, ψ and ϕ (as well as their anti-particles) produced in the decay of the Majorana fields N_i will sequentially decay, with ψ (and its anti-particle) decaying into the visible sector fields and ϕ (and its anti-particle) decaying into the dark sector fields. Their decays thus produce a net lepton asymmetry in the visible sector and a lepton asymmetry of opposite sign in the dark sector. We note that the absence of the decays $\psi \to \bar{X} + X'$ and $\phi^* \to L + H$ guarantees that leptonic asymmetries of equal and opposite sign are generated in the visible and in the dark sectors. Indeed, right after the heavy Majorana fermions N_i have decayed completely, and created the excess of ψ, ϕ over $\bar{\psi}, \phi^*$, equal and opposite lepton numbers are already assigned to the visible sector and the dark sector. It is clear from the above analysis that there is no violation of lepton number in the entire process of generating the leptonic asymmetries. We further note that while sphaleron interactions are active during the period when the leptogenesis and the genesis of (asymmetric) dark matter occur, they are not responsible for creating a net $B–L$ number in the visible sector, though they do play a role in transmuting a part of the lepton number into baryon number in the visible sector.

As will be discussed in Section 4, the symmetric component of dark matter would be sufficiently depleted by annihilating via a Z' pole into standard model particles, which ensures the asymmetric dark matter to be the dominant component of the current dark matter relic abundance. One can estimate on general grounds the mass of the dark particles in this model for the cosmic coincidence to occur. Since the total $B–L$ in the universe vanishes, the $B–L$ number in the visible sector is equal in magnitude and opposite in sign to the lepton number created in the visible sector right after

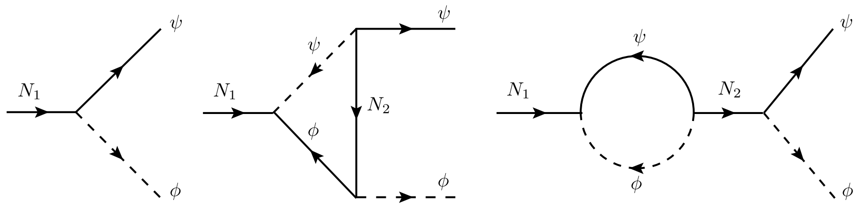

Fig. 1. An exhibition of the generation of asymmetry in ψ, ϕ over their anti-particles $\bar{\psi}, \phi^*$ from the decay of the Majorana field N_1. The lepton number is conserved in these processes.

N_i have completely decayed (the decay of N_i does not generate any baryon asymmetry), and thus is equal to the lepton number in the dark sector, i.e.,

$$(B–L)_v = L_d, \tag{5}$$

where the indices v, d denote the visible sector and the dark sector respectively. We are interested in the relative density of particle species at the time when the sphaleron interactions go out of the thermal equilibrium. This happens at a temperature of ~ 100 GeV which lies below the top mass so that the top quark would have already decoupled and no longer participates in the thermal bath. After the decoupling of the sphaleron interactions B and L are separately conserved and correspond to the B and L seen today. An analysis of the chemical potentials [17,5] allows us to compute the current value of B in term of $(B–L)_v$ so that

$$\frac{B^f}{(B–L)_v} = \frac{30}{97}, \tag{6}$$

where B^f denotes the final (and currently observed) value of the baryon number density. Assuming that X and X' have the same mass, and using Eqs. (2), (5) and (6) we obtain the mass of the dark particles

$$m_X = m_{X'} \approx 0.85 \text{ GeV}. \tag{7}$$

We turn now to the detail of the generation of the asymmetry between ψ, ϕ and $\bar{\psi}, \phi^*$. We assume there are two Majorana fields N_1 and N_2 with N_2 mass M_2 being much larger than the N_1 mass M_1, i.e., $M_2 \gg M_1$. Thus the generation of the asymmetry is mostly through the decay of N_1. The diagrams that contribute to it are shown in Fig. 1 where the Majorana particles N_i decay into the Dirac fermion ψ and the complex scalar ϕ with ψ and ϕ carrying opposite lepton numbers while the Majorana fields N_i carry no lepton number. As is well-known one needs an interference of the tree and the loop diagrams to create the asymmetry. The loop diagrams consist of a vertex diagram and a wave function diagram as shown in Fig. 1. The excess of ψ, ϕ over $\bar{\psi}, \phi^*$ is given by[3]

$$\epsilon = \frac{\Gamma(N_1 \to \psi\phi) - \Gamma(N_1 \to \bar{\psi}\phi^*)}{\Gamma(N_1 \to \psi\phi) + \Gamma(N_1 \to \bar{\psi}\phi^*)} \simeq -\frac{1}{8\pi}\frac{\mathrm{Im}(\lambda_1^2\lambda_2^{*2})}{|\lambda_1|^2}\frac{M_1}{M_2}, \tag{8}$$

where we have included both the vertex contribution and the wave contribution. Since the dark sector does not communicate with the visible sector, $(B–L)_v$ is equal in magnitude and opposite in sign to the lepton number generated in the visible sector,

Table 1
Lepton numbers and $U(1)_x$ charges of the superfields that enter in the generation of leptonic asymmetries for a gauged $U(1)_x$ model.

	$\hat{N}_i$	$\hat{Y}$	$\hat{Y}'$	$\hat{X}$	$\hat{X}^c$	$\hat{X}'$	$\hat{X}'^c$
L	0	−1	+1	$-\frac{1}{2}$	$+\frac{1}{2}$	$-\frac{1}{2}$	$+\frac{1}{2}$
$U(1)_x$	0	0	0	+1	−1	−1	+1

$$(B–L)_v = -L_v = -\frac{3}{4}\frac{\kappa\epsilon\zeta(3)g_N T^3}{\pi^2}, \tag{9}$$

where $\zeta(3) \sim 1.202$, $g_N = 2$ for the Majorana field N_1, κ is the washout factor [18] due to inverse processes $\psi + \phi \to N_1$, $\bar{\psi} + \phi^* \to N_1$ and we assume $\kappa \sim 0.1$. Using Eq. (6), one could further link the current baryon number to the excess of ψ, ϕ over $\bar{\psi}, \phi^*$ as

$$\frac{B^f}{s} = \frac{30}{97}\frac{(B–L)_v}{s} = -\frac{30}{97}\frac{135\zeta(3)}{4\pi^4}\frac{\kappa\epsilon}{g_s}, \tag{10}$$

where the entropy density $s = 2\pi^2 g_s T^3/45$ and $g_s \approx 100$ is the entropy degrees of freedom at $T \sim 100$ GeV when the sphaleron interactions decouple. Using the current astrophysical constraint given in Eq. (1) and Eq. (10) we estimate $|\epsilon| \sim 10^{-6}$.

3. Supersymmetric model

For the supersymmetric case we choose the following set of fields: $(\hat{N}_i(i \geqslant 2), \hat{Y}, \hat{Y}', \hat{X}, \hat{X}^c, \hat{X}', \hat{X}'^c)$ where ˆ denotes superfields, and their lepton numbers and $U(1)_x$ charges are summarized in Table 1. From the table it is clear that $U(1)_x$ is anomaly free and can be gauged.

For these superfields we assume a superpotential of the following form which conserve both the lepton number and the $U(1)_x$ gauge symmetry:

$$W = W_Y + W_m, \tag{11}$$

where W_Y contains the Yukawa couplings

$$W_Y = \lambda_i \hat{N}_i \hat{Y}\hat{Y}' + \beta\hat{Y}\hat{L}\hat{H} + \beta'\hat{Y}\hat{X}^c\hat{X}'^c + \gamma\hat{Y}'\hat{X}\hat{X}', \tag{12}$$

and W_m contains the mass terms

$$W_m = M_i\hat{N}_i\hat{N}_i + M_Y\hat{Y}\hat{Y}' + m_X\hat{X}\hat{X}^c + m_{X'}\hat{X}'\hat{X}'^c. \tag{13}$$

For the supersymmetric model, a possible candidate for $U(1)_x$ is $U(1)_{B-L}$ if one includes three right-handed neutrinos to the particle spectrum. Along with the anomaly free spectrum of Table 1, one can then gauge $U(1)_{B-L}$. In this case $\hat{Y}, \hat{Y}'$ along with the dark matter fields $\hat{X}, \hat{X}^c, \hat{X}', \hat{X}'^c$ will all carry $U(1)_{B-L}$ charges as shown in Table 2. And of course, the minimal supersymmetric standard model matter fields also carry $U(1)_{B-L}$ quantum numbers. In this case we require that all the fundamental interactions conserve the lepton number and the $U(1)_{B-L}$ gauge symmetry, and the superpotentials of Eqs. (12) and (13) remain unchanged. This model has the very interesting feature in that we have a gauged $U(1)_{B-L}$

[3] This calculation is similar to the calculation done in [12] for leptogenesis where the heavy Majorana fields N_i decay to L, H and their anti-particles. The difference here is that for leptogenesis, the wave contribution has two diagrams due to L, H being $SU(2)$ doublets; whereas for our case there is only one diagram for the wave contribution.

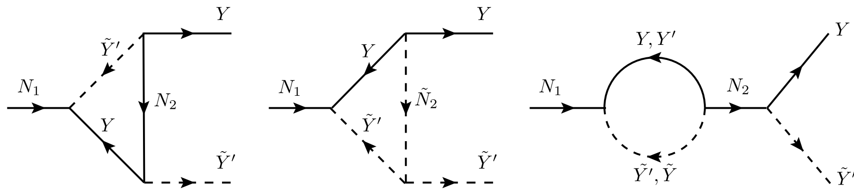

Fig. 2. Loop diagrams which are responsible for the genesis of asymmetry from the decay of N_1 to $Y\tilde{Y}'$. There are similar diagrams for the decay of N_1 to $\tilde{Y}Y'$, and for the decay of $\tilde{N}_1$ to YY' and to $\tilde{Y}\tilde{Y}'$. The lepton number is conserved in these processes.

Table 2
Lepton numbers and $U(1)_{B-L}$ charges of the superfields that enter in the generation of leptonic asymmetries for a gauged $U(1)_{B-L}$ model.

	$\hat{N}_i$	$\hat{Y}$	$\hat{Y}'$	$\hat{X}$	$\hat{X}^c$	$\hat{X}'$	$\hat{X}'^c$
L	0	-1	$+1$	$-\frac{1}{2}$	$+\frac{1}{2}$	$-\frac{1}{2}$	$+\frac{1}{2}$
$U(1)_{B-L}$	0	$+1$	-1	$+\frac{1}{2}$	$-\frac{1}{2}$	$+\frac{1}{2}$	$-\frac{1}{2}$

which leads to a conserved B–L with the initial condition B–$L=0$ in the universe.

All the following discussions in this section apply to both of the above two models. As in the non-supersymmetric case, for the generation of the asymmetry, we assume λ_i to be complex, and β, β', γ are assumed to be real.[4] Again as in the non-supersymmetric case we assume $i=2$ and assume the $\hat{N}_2$ mass M_2 to be much larger than the $\hat{N}_1$ mass M_1, and $M_Y \gg m_X + m_{X'}$. Again, M_Y could lie in a broad range from order of TeV to much higher scales. From the interactions of Eq. (12) we see that $\hat{Y}'$ decays exclusively into the dark sector so that $\hat{Y}' \to \hat{X} + \hat{X}'$ while $\hat{Y}$ could decays into the visible sector as well as dark sector particles. However, with the assumption $|\beta'| \ll |\beta|$, $\hat{Y}$ will decay dominantly into visible sector particles, i.e., $\hat{Y} \to \hat{L} + \hat{H}$.

As in the non-supersymmetric case the asymmetries in $\hat{Y}$ and in $\hat{Y}'$ are generated via the interference of the tree level amplitudes with the loop diagrams as shown in Fig. 2. The excess of $\hat{Y}, \hat{Y}'$ over their anti-particles $\bar{\hat{Y}}, \bar{\hat{Y}}'$ is given by a sum of several ϵ's, where these ϵ's are defined by the final decaying products. For example, one of these ϵ's is defined by

$$\epsilon_{Y\tilde{Y}'} = \frac{\Gamma(N_1 \to Y\tilde{Y}') - \Gamma(N_1 \to \bar{Y}\tilde{Y}'^*)}{\Gamma(N_1 \to Y\tilde{Y}') + \Gamma(N_1 \to \bar{Y}\tilde{Y}'^*)}. \tag{14}$$

Similarly one could define $\epsilon_{\tilde{Y}Y'}$ to parameterize the excess of $\tilde{Y}Y'$ over $\tilde{Y}^*\bar{Y}'$ decays from N_1; $\epsilon_{YY'}$ for the excess of YY' over $\bar{Y}\bar{Y}'$ decays from $\tilde{N}_1$; and $\epsilon_{\tilde{Y}\tilde{Y}'}$ for the excess of $\tilde{Y}\tilde{Y}'$ over $\tilde{Y}^*\tilde{Y}'^*$ decays from $\tilde{N}_1$. Similar to the non-supersymmetric case the total asymmetry is a sum of the asymmetries arising from the interference of the tree diagram with the vertex diagrams and with the wave function diagram. An analysis [12,8] of the asymmetries gives the following relation

$$\epsilon_{Y\tilde{Y}'} = \epsilon_{\tilde{Y}Y'} = \epsilon_{YY'} = \epsilon_{\tilde{Y}\tilde{Y}'} \equiv \varepsilon, \tag{15}$$

and in the limit $M_2 \gg M_1$ we obtain [8]

$$\varepsilon \simeq -\frac{1}{4\pi}\frac{\mathrm{Im}(\lambda_1^2\lambda_2^{*2})}{|\lambda_1|^2}\frac{M_1}{M_2}. \tag{16}$$

The difference between the front factor in Eq. (16) and the front factor in Eq. (8) is due to the fact that there are two vertex diagrams for the supersymmetric case (see Fig. 2) compared to just one vertex diagram for the non-supersymmetric case (see Fig. 1). Thus the total excess of $\hat{Y}$ over $\bar{\hat{Y}}$, i.e., $Y, \tilde{Y}$ over $\bar{Y}, \tilde{Y}^*$ generated by the decay of $\hat{N}_1$ is given by:

$$\Delta n_Y \equiv n_{\hat{Y}} - n_{\bar{\hat{Y}}}, \tag{17}$$

where Δn_Y is computed to be

$$\Delta n_Y = \left[\frac{3}{4}(\epsilon_{Y\tilde{Y}'} + \epsilon_{\tilde{Y}Y'}) + (\epsilon_{YY'} + \epsilon_{\tilde{Y}\tilde{Y}'})\right]\frac{\kappa\zeta(3)g_N T^3}{\pi^2}.$$

Here the factor of $\frac{3}{4}$ is for N_1 and a factor of 1 for $\tilde{N}_1$, and again κ is a washout factor which we assume to be 0.1. The excess of $\hat{Y}, \hat{Y}'$ then gives rise to an equal but opposite lepton number to the visible sector and to the dark sector. Thus we obtain the (B–L)-number density in the visible sector to be

$$(B\text{–}L)_v \approx 2\kappa\varepsilon/g_s, \tag{18}$$

where again $g_s \approx 100$ is the entropy degrees of freedom at $T \sim 100$ GeV when the sphaleron interactions decouple. Similar to the discussion in the non-supersymmetric case, we estimate $|\varepsilon| \sim 10^{-6}$.

The analysis of the dark matter mass in the supersymmetric model is also very similar to the one in the non-supersymmetric case, and Eq. (5) still holds. The computation of B^f will be identical to the non-supersymmetric case and Eq. (6) also holds. This is so because the sleptons and squarks have already decayed into standard model particles and the memory of them is lost by the time sphaleron interactions go out of thermal equilibrium after which B and L are separately conserved. The modification that will occur is due to the presence of additional fields $\hat{X}^c, \hat{X}'^c$ and both the bosonic and fermionic components of the superfields should be considered in the analysis. However, the total lepton number will not be affected by the number of fields. Assuming the bosonic and the fermionic fields $\hat{X}, \hat{X}', \hat{X}^c, \hat{X}'^c$ all have the same mass, and again using Eqs. (2), (5) and (6) we obtain

$$m_X = m_{X'} = 0.85\ \text{GeV}. \tag{19}$$

The $U(1)_x$ gaugino λ_x is given a soft mass $\mathcal{L}_{\lambda_x} = -m_\lambda\bar{\lambda}_x\lambda_x$. Assuming $m_\lambda > m_X + m_{\tilde{X}}$, the gaugino λ_x can decay into $X\tilde{X}$ or $X'\tilde{X}'$, etc, via the supersymmetric interaction $\mathcal{L} \sim \bar{\lambda}_x X\tilde{X} + \bar{\lambda}_x X'\tilde{X}' + \bar{\lambda}_x X^c\tilde{X}^c + \bar{\lambda}_x X'^c\tilde{X}'^c + h.c.$ Thus the gaugino λ_x decays into dark particles and is removed from the low energy spectrum.

4. Phenomenology

We discuss now phenomenological implications of the model. An interesting implication of our model arises in the neutrino sector. Here we add three families of right-handed neutrinos. Now we also assume the coupling β is family dependent, i.e., $\beta \to \beta_i$ where $i = 1, 2, 3$ correspond to e, μ, τ, cf., Eqs. (3) and (12). The terms which will contribute to neutrino masses read

$$\mathcal{L}_m = \beta_i\bar{\psi}_R L_i H + \beta''_{ij}\bar{\nu}_{iR}L_j H + \mu'_i\bar{\nu}_{iR}\psi_L + h.c. \tag{20}$$

[4] The interactions $\hat{N}_i\hat{X}\hat{X}^c$ and $\hat{N}_i\hat{X}'\hat{X}'^c$ could exist. However, the inclusion of these two interactions will not change our discussion. This is so because $(\hat{X}, \hat{X}^c)$ and $(\hat{X}', \hat{X}'^c)$ carry opposite lepton numbers, and thus there will be no net lepton number generated in the dark sector through $\hat{N}_i$ decay from these interactions. Here we assume $\hat{N}_i$ would mostly decay into $\hat{Y}, \hat{Y}'$.

After spontaneous breaking of the electroweak symmetry, the full mass terms recast into

$$\mathcal{L}_m = \vec{\nu}_R^T \mathcal{M} \vec{\nu}_L + h.c., \tag{21}$$

where we have defined

$$\vec{\nu}_R^T = \left(\bar{\nu}_R^e, \bar{\nu}_R^\mu, \bar{\nu}_R^\tau, \bar{\psi}_R\right), \tag{22}$$

$$\vec{\nu}_L^T = \left(\nu_L^e, \nu_L^\mu, \nu_L^\tau, \psi_L\right), \tag{23}$$

$$\mathcal{M} = \begin{pmatrix} m_{ee}^\nu & m_{e\mu}^\nu & m_{e\tau}^\nu & \mu'_1 \\ m_{e\mu}^\nu & m_{\mu\mu}^\nu & m_{\mu\tau}^\nu & \mu'_2 \\ m_{e\tau}^\nu & m_{\mu\tau}^\nu & m_{\tau\tau}^\nu & \mu'_3 \\ \mu_1 & \mu_2 & \mu_3 & m_1 \end{pmatrix}, \tag{24}$$

and

$$\mu_i = \frac{1}{\sqrt{2}}\beta_i v, \qquad m_{ij}^\nu = \frac{1}{\sqrt{2}}\beta''_{ij} v, \tag{25}$$

and where v is the VEV of Higgs. In matrix $\mathcal{M}$, m_1 is much larger than all the other entries.

A diagonalization of matrix $\mathcal{M}$ gives four Dirac fermions in the mass eigenbasis: three of which correspond to the three neutrinos, while the fourth one is mostly constituted by ψ which is much heavier. However, a fine-tuning is needed to get the light neutrino masses in the experimental range. $\mathcal{M}$ can be diagonalized by using a biunitary transformation so that

$$V^\dagger \mathcal{M} U = \mathcal{M}_D. \tag{26}$$

Thus the left-handed neutrino states transform as

$$\nu_{iL} = \sum_{a=1}^{4} U_{ia} \nu'_{aL}, \tag{27}$$

where ν'_{aL} are in the mass diagonal basis. Eq. (27) implies that, for example, the partial decay widths of the W and Z bosons will be modified so that

$$\Gamma(W \to \ell_i \bar{\nu}_i) = \Gamma(W \to \ell_i \bar{\nu}_i)_{\rm SM}\left(1 - |U_{i4}|^2\right), \tag{28}$$

$$\Gamma(Z \to \nu_i \bar{\nu}_i) = \Gamma(Z \to \nu_i \bar{\nu}_i)_{\rm SM}\left(1 - |U_{i4}|^2\right)^2. \tag{29}$$

Now the low energy electroweak data is in excellent agreement with the standard model and thus the new physics can be accommodated only within the error bars. Here we use the data on the hidden decays of the Z boson [19], which in the standard model are neutrinos, to constrain U_{i4}, i.e.,

$$\Gamma(Z \to \nu \bar{\nu}) = (499 \pm 1.5)\ \text{MeV}. \tag{30}$$

Using Eq. (30) and assuming the correction U_{i4} is uniform across generations we get an upper limit on U_{i4} of

$$|U_{i4}| \lesssim 4 \times 10^{-2}. \tag{31}$$

The presence of a sizable U_{i4} will also affect other electroweak processes where neutrinos appear. Thus more accurate measurements in the electroweak sector in the future, for example, at the ILC could reveal the presence of a non-negligible value of U_{i4}. This would provide a possible test of the model.

Next we demonstrate that a sizable U_{i4} can be obtained from Eq. (21) consistent with small neutrino masses. We first consider an example of one generation of neutrino (say the third generation) mixing with the ψ field. For this case we have

$$\mathcal{L}_m^{(3)} = (\bar{\nu}_{3R}, \bar{\psi}_R)\begin{pmatrix} m_{\tau\tau}^\nu & \mu'_3 \\ \mu_3 & m_1 \end{pmatrix}\begin{pmatrix} \nu_{3L} \\ \psi_L \end{pmatrix} + h.c. \tag{32}$$

With the inputs $m_{\tau\tau} = 10^{-12}$, $\mu'_3 = 10^{-9}$, $\mu_3 = 10$, $m_1 = 1000$ (all masses in GeV), we obtain the mass eigenvalue of the neutrino to be around 10^{-2} eV, $U_{34} \sim 0.01$ consistent with the constraint of Eq. (31). The 2×2 matrix analysis above uses a lopsided matrix in Eq. (32). An analysis of the lopsided 4×4 case is more elaborate and for that reason we do not give an extended analysis of this case here but we expect that a sizable U_{i4} can be generated in that case as well.

We discuss now another sector of the parameter space of Eq. (21). Here we assume a symmetrical form for the neutrino mass terms so that

$$\mathcal{L}_m^\nu = \vec{\nu}_R^T \begin{pmatrix} m_{\nu_e} & 0 & 0 & \mu'_1 \\ 0 & m_{\nu_\mu} & 0 & \mu'_2 \\ 0 & 0 & m_{\nu_\tau} & \mu'_3 \\ \mu_1 & \mu_2 & \mu_3 & m_1 \end{pmatrix} \vec{\nu}_L + h.c., \tag{33}$$

and we further assume

$$\mu_1 = \mu'_1, \qquad \mu_2 = \mu'_2, \qquad \mu_3 = \mu'_3. \tag{34}$$

The matrix of Eq. (33) contains no direct mixings among the neutrino flavor states. However, we will see that their mixings with the field ψ automatically lead us to neutrino flavor mixings. To exhibit this result we diagonalize the matrix of Eq. (33) by an orthogonal transformation.

By setting $m_{\nu_e} = 10^{-11}$, $m_{\nu_\mu} = 1.7 \times 10^{-10}$, $m_{\nu_\tau} = 2 \times 10^{-9}$, $m_1 = 2000$, $\mu_1 = 3.6 \times 10^{-5}$, $\mu_2 = 8.9 \times 10^{-5}$, $\mu_3 = 5.9 \times 10^{-4}$ (all masses in GeV) the three neutrino masses in the mass diagonal basis are calculated to be

$$m_3 \approx 4.8 \times 10^{-2}\ \text{eV}, \tag{35}$$

$$m_2 \approx 1.2 \times 10^{-2}\ \text{eV}, \tag{36}$$

$$m_1 \approx 4.2 \times 10^{-3}\ \text{eV}, \tag{37}$$

which produce the normal hierarchy of neutrino masses [19] and the mass eigenvalue of the heavy field ψ is still approximately m_1. For the neutrino mixings we obtain

$$\sin^2\theta_{12} \approx 0.30, \qquad \sin^2\theta_{23} \approx 0.36, \qquad \sin^2\theta_{13} \approx 0.024, \tag{38}$$

while the experimental values are [19]

$$\sin^2\theta_{12} = 0.307^{+0.018}_{-0.016}, \qquad \sin^2\theta_{23} = 0.386^{+0.024}_{-0.021},$$
$$\sin^2\theta_{13} = 0.0244^{+0.0023}_{-0.0025}. \tag{39}$$

We see that our analysis of Eq. (38) is in good accord with the experimental determination of the mixing angles as given in Eq. (39). Specifically the model is consistent with the result from the Daya Bay neutrino reactor experiment [20] of $\theta_{13} \sim 9°$. Thus it is very interesting that the model provides an explanation of the neutrino mixings at a fundamental level, in that the neutrino mixings arise as a consequence of the interaction of the neutrinos with the primordial Dirac field ψ which enters in leptogenesis which points to the cosmological origin of neutrino mixings.

Other implications of the model involve flavor changing processes. For the supersymmetric model of Eq. (12), after spontaneous breaking one has interactions of the charged Higgs H^+ with charged leptons and Y:

$$\mathcal{L}_{H\ell\psi} = \beta_i \bar{Y} \ell_i H^+ + h.c., \tag{40}$$

where ℓ_i denotes the charged leptons. Such interactions will give rise to $\ell_i \to \ell_j \gamma$ processes, where a charged lepton ℓ_i converts into a charged lepton ℓ_j via exchange of Y while a photon is emitted by the charged Higgs inside the loop, see Fig. 3. Assuming $m_Y^2 \gg m_{H^+}^2$,

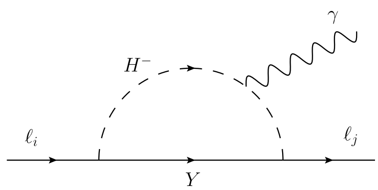

Fig. 3. Flavor changing processes $\ell_i \to \ell_j \gamma$ via the charged Higgs and Y loop.

we obtain the decay rate of the flavor changing process $\ell_i \to \ell_j \gamma$ to be[5]

$$d\Gamma_{\ell_i \to \ell_j \gamma} = \frac{\alpha_{\rm em}(\beta_i \beta_j)^2}{(16\pi^2)^2} \frac{m_i^3}{M_Y^2}, \tag{41}$$

where m_i is the mass of the decaying charged lepton and we have used $m_i \gg m_j$. The current experimental upper bounds on the branching ratio of such flavor changing processes read [21,22]

$$\mathcal{B}r(\mu \to e\gamma) \lesssim 2.4 \times 10^{-12}, \tag{42}$$

$$\mathcal{B}r(\tau \to e\gamma) \lesssim 3.3 \times 10^{-8}, \tag{43}$$

$$\mathcal{B}r(\tau \to \mu\gamma) \lesssim 4.4 \times 10^{-8}. \tag{44}$$

Using the mean lifetimes for μ and τ [19], the branching ratios Eqs. (42)–(44) and Eq. (41), and $M_Y \sim 1$ TeV we obtain

$$\beta_1 \sim \beta_2 \lesssim 3 \times 10^{-3}. \tag{45}$$

Once β_1, β_2 are fixed, one can estimate β_3 by

$$\beta_3 \lesssim 2 \times 10^{-4}/\beta_1. \tag{46}$$

One can expect observable effects in these flavor changing processes in future experiments with improved sensitivities. And at the same time, we see that with these constraints, μ_3 could be of $\mathcal{O}(10)$ GeV, cf., Eq. (25), thus one would also expect to see the effect we discussed at Eq. (32).

Next we discuss the phenomenological implications of the model in the dark sector. An important issue concerns the dissipation of thermally produced dark matter. To dissipate the symmetric component of dark matter we use the fact that dark matter is charged under the gauged group $U(1)_x$ or $U(1)_{B-L}$. We assume that the $U(1)$ gauge boson gains mass via the Stueckelberg mechanism [23].[6] For the gauged $U(1)_x$ model, one could assume a kinetic mixing of the $U(1)_x$ gauge boson with the $U(1)_Y$ gauge boson [25]. This mechanism allows one to dissipate the symmetric component of dark matter which can annihilate into the standard model particles via the $Z - Z'$ mixing. The analysis here is very similar to the ones discussed in [8].

The Z' gauge boson can make a contribution to the anomalous magnetic moment of the muon. At the one loop order one finds

$$\Delta a_\mu \simeq \delta^2 \frac{g_Y^2 m_\mu^2}{48\pi^2 M_{Z'}^2}, \tag{47}$$

where δ is the coupling of the Z' with matter current, i.e., $\mathcal{L}_{\rm int}^{Z'} = \delta Z'_\mu J^\mu$. A value of $\delta \sim 10^{-3}$ and a Z' mass of order of a few GeV gives Δa_μ significantly below the current experimental limit on the deviation of a_μ from the standard model value of $\Delta(a_\mu) < 3 \times 10^{-9}$ [19]. At the same time δ and $M_{Z'}$ satisfy the LEP II constraint of [26] that $M_{Z'}/g_{Z'\bar{\ell}\ell} > 6$ TeV.

As discussed earlier, fields X, X' (non-supersymmetric case) or $\hat{X}, \hat{X}^c, \hat{X}', \hat{X}'^c$ (supersymmetric case) constitute the dark matter which are all light with masses $\mathcal{O}(1)$ GeV. Since the coupling between Z' and standard model particles can only be $\sim 10^{-3}$ because of experimental constraints, a sufficient depletion of the symmetric component of dark matter (up to or less than 10% of the total dark matter relic density), requires a Breit–Wigner enhancement, so that the Z' mass is around twice the dark matter mass. It is seen that with a kinetic mixing parameter $\delta \sim 0.001$ [27], for a dark matter mass of ~ 1 GeV, a Z' mass of ~ 3 GeV does allow the symmetric component of dark matter to be depleted down to less than 10% of the total dark matter relic density. Thus the current dark matter would be constituted of up to 90% or more of the light asymmetric dark matter.

Such dark matter can scatter from quarks within a nucleon through the t-channel exchange of the Z' boson. The spin-independent dark matter-nucleon (target-independent) cross section can be approximately written as [28,29]

$$\sigma_{\rm SI} \sim \frac{4}{\pi} \frac{\delta^2 g_x^2 g_Y^2 \cos^4\theta_W \mu_n^2}{m_{Z'}^4}, \tag{48}$$

where μ_n is the dark matter-nucleon reduced mass. Using the parameters discussed above we find $\sigma_{\rm SI} \sim 10^{-37}$ cm^2, which is just on the edge of sensitivity of the CRESST I experiment [30]. Thus improved experiments in the future in the low dark matter mass region with better sensitivities should be able to test the model.

For the supersymmetric gauged $U(1)_{B-L}$ model, without using the kinetic mixing mechanism, one could use the $U(1)_{B-L}$ gauge boson to dissipate the symmetric component of dark matter. As discussed in [5], the mass of the $U(1)_{B-L}$ gauge boson can lie in a few GeV range and be consistent with the LEP II constraints and with the UA2 cross section bounds [31]. The analysis of [5] also shows that the symmetric component of dark matter can be sufficiently depleted.

There could be also indirect hints for the existence of the asymmetric dark matter. For example, assume that dark matter consists of both an asymmetric component which is the dominant one ($\gtrsim 90\%$) and a subdominant component ($\lesssim 10\%$) which is WIMP like. A detailed analysis shows that the subdominant component could still be detected [5]. On the other hand the WIMP model would not constitute the entire relic density which would require the asymmetric dark matter to make up the deficit. This could provide an indirect evidence for asymmetric dark matter if WIMPs were observed in direct detection but a detailed theory model shows a large deficit in its contribution to the relic density.

Thus quite interestingly the above discussion indicates that the cogenesis model which relates to cosmological issue gets directly related to particle physics experiments specifically experiments at the intensity frontier [32] and those related to search for dark matter.

Finally, we compare our work briefly with the work of other authors specifically the works of [13,14]. There are major differences between our work and theirs both in the structure of the model as well as in the phenomenological implications. At the level of the structure of the model the major difference between our model and the models of [13,14] is that for our model the asymmetries in both the visible and the dark sectors are generated through the decay of heavy Majorana fields, which do not carry any lepton or

[5] Loops which involve the Higgsinos and $\tilde{Y}$ also contribute to $\ell_i \to \ell_j \gamma$ process. A computation shows that these loops are suppressed by a factor of $m_{\tilde{H}}^2/m_{\tilde{Y}}^2$ compare to the charged Higgs and Y loop. For the case $m_{\tilde{Y}}^2 \gg m_{\tilde{H}}^2$, we could omit these contributions.

[6] An alternative way of depleting the symmetric component of the dark matter is assuming the $U(1)_x$ gauge boson to be massless (dark photon). Then the symmetric component of dark matter could sufficiently annihilate into the $U(1)_x$ dark photons and become radiation in the early universe. As shown in [24], the constraints on the number of extra effective neutrino species, $\Delta N_{\rm eff}$, can be satisfied for a large class of asymmetric dark matter models.

baryon numbers, while for the model of [13] the asymmetries in both sectors arise from the decay of heavy Dirac particles which carry baryon number and for the model of [14] the asymmetries arise from the decay of heavy complex scalars which carry either baryon or lepton number.

Further, in our model we have two mediator fields (ψ, ϕ in the non-supersymmetric case and $\hat{Y}, \hat{Y}'$ in the supersymmetric case) which subsequently decay into visible or dark sector particles after they are produced by the decay of the heavy Majorana fields, and this procedure has the advantage that the experimental data on the asymmetry in the universe does not set a bound on the mass of the mediator fields or on the couplings of the mediator fields to the standard model particles. In addition, our model is focused on generating the asymmetry in the leptonic sector, whereas the work of [13] focused on generating the asymmetry in the baryonic sector. Although the work of [14] also has a model on generating the asymmetry in the leptonic sector, that model is very different from ours.

In addition to differences in the theoretical structure of the models, there are very significant phenomenological differences between the model presented here and the works of [13,14]. The phenomenological implications of the works of [13,14] have been spelled out in these works and we do not wish to enumerate them here. One item, however which we wish to point out it that in the model of [13] that dark particles can induce proton decay. This feature is not shared by our model. In contrast we have discussed in this work a variety of phenomena arise from the leptonic sector, which can provide low energy tests of the proposed model. These include implication of the model for electroweak physics, neutrino masses and mixings and lepton flavor changing processes. Further, the mechanism for the dissipation of symmetric component of dark matter in the model is also different. In particular, we propose that the symmetric component of dark matter can annihilate efficiently into standard model particles through a very light Z' with mass around twice the dark matter mass [5,8]. This Z' gauge boson can only couple very weakly to the standard model particles and thus satisfies all the current experimental constraints. In summary, both the theoretical structure and the phenomenological implications of the proposed models are very different from previous works on this topic.

5. Conclusion

In this work we discussed models of leptogenesis where the fundamental interactions do not violate lepton number, and the total B–L in the universe vanishes. Thus the generation of a net lepton number in the visible universe is compensated by the generation of an equal amount of anti-lepton number in the dark sector. Baryogenesis in this class of models occurs via the sphaleron interactions which convert a part of the lepton number in the visible sector into baryon number in the visible sector. Three models are discussed in this work: one non-supersymmetric gauged $U(1)_x$ model, one supersymmetric gauged $U(1)_x$ model, and another supersymmetric gauged $U(1)_{B-L}$ model. A detailed analysis shows that the models can generate the baryon number density to the entropy density ratio consistent with the observed value. The models also produce the desired amount of dark matter and provide an explanation of the observed dark matter density to the baryonic matter density in the universe. Thus the proposed models provide a possible explanation of the three cosmology puzzles mentioned in the introduction. The results from the analysis of these models indicate that a violation of lepton number in the fundamental interactions is not essential for leptogenesis. The symmetric component of dark matter in these models is dissipated via kinetic mixing between $U(1)_x$ and $U(1)_Y$ gauge bosons, or for the supersymmetric gauged $U(1)_{B-L}$ model through a light Z'_{B-L} gauge boson. The gauged $U(1)_{B-L}$ model is rather attractive in that the $U(1)_{B-L}$ gauge invariance requires conservation of B–L and B–$L = 0$ provides the most natural initial condition for the universe.

Phenomenological implications of the models were discussed. These include implications for electroweak physics and neutrino masses and mixings. Specifically it is seen that small corrections arise in the electroweak sector which may be detectable in future high precision machines such as ILC. Further, it is seen that in the proposed models neutrino interactions with a primordial Dirac field which enters in leptogenesis naturally lead to neutrino mixings after spontaneous breaking of the electroweak symmetry. For the supersymmetric case, our model may lead to the charged lepton flavor changing processes such as $\ell_i \to \ell_j \gamma$ at the loop level. Also discussed is the feasibility of the direct detection of dark matter in $\mathcal{O}(1)$ GeV mass range. The models proposed in this work are significantly different from other cogenesis models both in theoretical structure as well as in their phenomenological implications.

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