Leptogenesis as a Common Origin for Matter and Dark Matter

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Abstract

We propose a model of asymmetric dark matter (DM) where the dark sector is an identical copy of both forces and matter of the standard model (SM) as in the mirror universe models discussed in literature. In addition to being connected by gravity, the SM and DM sectors are also connected at high temperature by a common set of heavy right-handed Majorana neutrinos via their Yukawa couplings to leptons and Higgs bosons. The lightest nucleon in the dark (mirror) sector is a candidate for dark matter. The out of equilibrium decay of right-handed neutrino produces equal lepton asymmetry in both sectors via resonant leptogenesis which then get converted to baryonic and dark baryonic matter. The dark baryon asymmetry due to higher dark nucleon masses leads to higher dark matter density compared to the familiar baryon density that is observed. The standard model neutrinos in this case acquire masses from the inverse seesaw mechanism. A kinetic mixing between the $U(1)$ gauge fields of the two sectors is introduced to guarantee the success of Big-Bang Nucleosynthesis.
I. INTRODUCTION

It now appears established that dark matter accounts for about one quarter of the energy density, $\Omega$, of the universe and plays an essential role in the formation of large scale structure in it. The identity of dark matter, however, remains unknown since all the particles in the successful standard model can be ruled out as candidates. What the dark matter particles are, how they interact with visible matter and how their relic abundance originates, constitute some of the fundamental mysteries of particle physics and cosmology today. Adding to this puzzle is the observation that baryon contribution to $\Omega$ is only about a fifth (i.e., of the same order) of the dark matter contribution. This raises the question: could the two have a common origin?

The most popular class of candidates for dark matter are the stable weakly interacting massive particles (WIMPs), which arise in many well-motivated TeV scale extensions of the standard model. Their stability is guaranteed by some symmetry. While the WIMP dark matter does not decay due to the symmetry, pairs of them can annihilate and their relic density is determined by the freeze-out of the annihilation from equilibrium to the SM particles. The fact that their observed relic density can be naturally explained by the weak scale annihilation cross section makes these models quite appealing. However in most models, the matter and dark matter contributions to $\Omega$ are unrelated.

Coming to the question of the symmetry that ensures the stability of WIMPs, in most models one uses a $Z_2$ symmetry (e.g., R-parity in supersymmetry or KK-parity in the case of extra dimension models). On the other hand, one could quite easily imagine that the stability of the WIMP is guaranteed by a continuous $U(1)$ symmetry similar to baryon number (call it “dark baryon number”). In such a case, observed dark matter density would represent an asymmetry between dark matter and anti-dark matter densities exactly as the case for the observed asymmetry between familiar matter and anti-matter. If these two asymmetries could arise from a common mechanism, it would be a major step towards understanding why their contributions to $\Omega$ are of the same order. It is the goal of this paper to propose such a scenario.

The baryon-anti-baryon asymmetry in the universe (BAU) is of order $10^{-10}$ and can arise from the laws of microphysics if the three conditions proposed by Sakharov [1] are satisfied, namely baryon number violation, CP violation and out-of-equilibrium in the early universe. The SM, however, fails to realize the third condition since it requires that the Higgs mass must be less than 40 GeV; even if this condition was satisfied, with three generations of fermions, it cannot explain why the asymmetry is as large as the one observed. An elegant way to generate the BAU is through leptogenesis [2] in the framework of seesaw mechanism [3] which naturally explains the smallness of the observed neutrino masses. In these models, lepton asymmetry is generated through the out-of-equilibrium decay of the very heavy right-handed (RH) neutrinos which then get converted to baryon asymmetry through the non-perturbative $B + L$ violating electroweak sphaleron process [4].

The appealing mechanism of baryogenesis via leptogenesis combined with the fact that the relic abundances of baryons and dark matter are of the same order of magnitude inspire us to think whether genesis of dark matter could also have its origin in a manner similar to leptogenesis. The mirror universe models discussed in the literature [5] appear to be a natural setting for this. In the mirror models, the universe has two kinds of matter and forces: one consisting of a standard model sector with forces and matter that we are familiar with, such as quarks, leptons, W, Z, etc., and a parallel sector which is an exact replica (e.g., consisting
of mirror duplicates of quarks and leptons, W, Z, γ, gluons and the Higgs boson) called the mirror sector \[5\]. Forces (except for gravity) in one sector do not affect the matter in the other sector. In what follows we will denote the mirror particles by a prime on a symbol. The two sectors communicate with each other by gravity and possibly some SM singlet interactions which are very weak at the current age of the Universe. There is a dark baryon and lepton number in the mirror sector which is the exact analog of the familiar baryon and lepton number. The main hypothesis of this paper is that the same leptogenesis mechanism that could be producing matter-anti-matter asymmetry, is also producing asymmetry of dark matter-anti-dark-matter. This then links the dark matter energy density of the Universe to that contributed by matter making them of the same order, providing a resolution of the puzzle stated in the beginning \[8\].

A key ingredient in our attempt to connect the matter asymmetry to dark matter asymmetry is the assumption that the visible and the mirror sectors talk to each other not only through gravity but also through a common set of three right-handed neutrinos coupled to leptons and Higgs fields in each sector through Yukawa couplings\[10\], as shown in Fig. 1. Since the RH neutrinos are standard model singlets, this is consistent with gauge invariance. Also mirror symmetry makes the $N^i H$ couplings on both sides equal. The out-of-equilibrium decays of right-handed neutrinos in the early universe can then produce lepton number asymmetries in both sectors, which are then transferred to baryon and dark baryon numbers through the sphaleron processes in each sector. If one imposes exact mirror symmetry on the theory, the primordial lepton asymmetries generated in each sector are equal and after sphaleron interaction produce the same number density for baryons and dark baryons in the early universe. Since we expect the symmetry breaking pattern in both sectors to be different for the model to be consistent with cosmology (see below), the resulting energy density contributions can be different and in the ratio $\Omega_B : \Omega_{DM} \approx 1 : 5$ if we require that mass of the dark baryons is five times the mass of the familiar SM baryons. This mass difference can arise from the difference in the scales of two $SU(3)$, strong interactions ($\Lambda_{QCD}, \Lambda'_{QCD}$). It turns out that this difference depends on the ratio of two electroweak scales $v_{wk}$ and $v'_{wk}$, with $v_{wk} : v'_{wk} \sim 1 : 10^3$ giving the required difference between $\Omega_B$ and $\Omega_{DM}$.

The spectra of the SM neutrinos $\nu_\alpha$ and the dark neutrinos $\nu'_\alpha$ are determined by the inverse seesaw \[11\] and type-I \[3\] seesaw mechanisms, respectively, with the mixing between the $\nu_\alpha$ and $\nu'_\alpha$ given by the ratio $v_{wk}/v'_{wk}$. The dark neutrinos can decay into the SM particles due to this mixing.

In addition to the common set right-handed neutrinos, the SM sector and the mirror sector can also be connected through Higgs interaction and the kinetic mixing between the $U(1)$ gauge bosons consistent with gauge invariance (as shown in Fig. II). The photon sector mixing is necessary for the model to be consistent with Big Bang Nucleosynthesis (BBN). In this work we assume the photon in the mirror sector acquires a mass around 50 MeV through the spontaneous symmetry breaking so that the mirror-electrically charged particles which are heavier, pair annihilates into it before the BBN, and the mirror photon itself decays to the electron-positron pair through the kinetic mixing. To generate this small mass two Higgs doublets are needed in the mirror sector. The kinetic mixing between the photon and mirror photon determines the lifetime of the mirror photon, so that success of BBN puts

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1 For earlier suggestion that mirror baryons constitute dark matter of the Universe, see \[9\]. Our model is, however, very different from these models in many respects.
a lower bound on this mixing. Furthermore, it also determines the interactions between the dark matter particles and the nucleons; therefore, using the current constraint from the direct detection experiments produces an upper bound on this mixing.

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The canonical leptogenesis [2] provides a connection with the baryon asymmetry with neutrino masses via the RH neutrinos needed for seesaw mechanism. In our model, since we

\[ L = -\lambda_{ij} N_j P_L (\ell_i H) - \xi_{ij} N_j P_L (\ell'_i H') - \frac{1}{2} N_j M_j N_j^c + \text{h.c.}, \]  

\[ \text{FIG. 1: The connection between the standard model and the dark sector.} \]
have a common set of RH neutrinos connecting to leptons in both sectors, the leptogenesis connects the lepton asymmetry in both sectors.

Leptogenesis could be of two types. For the case of hierarchical right-handed neutrino masses, $M_1 \ll M_2, M_3$, a population of $N_1$ can be thermally produced at temperature $T \sim M_1$ with negligible productions of $N_{2,3}$ and lepton asymmetry can be produced through out-of-equilibrium decay of $N_1$ providing the interactions are CP violating[6]. On the other hand, we could have at least two of the right-handed neutrinos highly degenerate, in which case we have resonant leptogenesis[7]. In the first case, the amount of lepton asymmetry is highly sensitive to the values of the leptonic Yukawa couplings whereas in the latter case, we could generate large lepton asymmetry regardless of the coupling values by choosing the degree of degeneracy between the RH neutrino masses. We will see below that within the constraints of our model, the alternative of resonant leptogenesis provides a more satisfactory framework. To see this, we first present the formula for lepton asymmetry:

$$\epsilon \equiv \frac{\sum_\alpha (\Gamma(N_1 \rightarrow H\ell_\alpha) - \Gamma(N_1 \rightarrow \bar{H}\bar{\ell_\alpha}))}{\sum_\alpha (\Gamma(N_1 \rightarrow H\ell_\alpha) + \Gamma(N_1 \rightarrow \bar{H}\bar{\ell_\alpha}))},$$

(2)

where the sum goes only over the SM sector fields and applies to both numerator and denominator. Correspondingly, $\epsilon'$ in the mirror sector can be defined in the same way with $H, \ell$ replaced by $H', \ell'$. The tree-level decay rates are given by

$$\Gamma_{H\ell} \approx \Gamma_{\bar{H}\bar{\ell}} \approx \frac{(\lambda^\dagger \lambda)_{11} M_1}{16\pi}, \quad \Gamma_{H'\ell'} \approx \Gamma_{\bar{H}'\bar{\ell}'} \approx \frac{(\xi^\dagger \xi)_{11} M_1}{16\pi},$$

(3)

therefore the total rate is $(\lambda^\dagger \lambda + \xi^\dagger \xi)_{11} M_1/(8\pi)$. The CP asymmetric parameters $\epsilon^{(l)}$ originate from the interference of tree-level and one-loop amplitudes. In our scenario, the self energy diagram gets two contributions, due to the fact that both $\ell, \ell'$ show up in the loop. For hierarchical right-handed neutrino spectrum, one obtains

$$\epsilon \approx \sum_{k \neq 1} \frac{1}{16\pi M_k} \text{Im}\left\{3[(\lambda^\dagger \lambda)_{k1}]^2 + 2[(\lambda^\dagger \lambda)_{k1}(\xi^\dagger \xi)_{k1}]\right\}/(\lambda^\dagger \lambda)_{11},$$

(4)

and $\epsilon' = \epsilon(\lambda \leftrightarrow \xi)$. For the resonant leptogenesis case, in the above expression, the factor $M_1/M_2$ is replaced by $M/\Gamma_N$ where $\Gamma$ is the decay width of the RH neutrino. The CP asymmetries first generate asymmetries in the (dark) lepton sectors. The evolution of this asymmetry can be studied using Boltzmann equations [6]. In our case, they differ from the conventional ones in the literatures since we have included CP asymmetries due to the on-shell part of the additional scattering processes. Other than this the discussion is standard [6] and we do not reproduce it here.

Above the (mirror) electroweak scales, the (mirror) lepton numbers will be transferred into a baryon asymmetry by the $SU(2)^{(l)}$ sphaleron processes. The baryon asymmetry can be written as

$$\eta_B \equiv \frac{n_B}{n_\gamma} \simeq \frac{3}{4} \cdot \frac{g^0_{ss}}{g_{ss}} \cdot \frac{28}{79} \cdot \kappa \cdot \epsilon,$$

(5)

where $g^0_{ss}$ and $g_{ss}$ are the total number of relativistic degrees of freedom today and during sphaleron, $28/79$ accounts for the sphaleron effects, $\kappa$ is the washout factor.
The resulting ratio of baryon to dark baryon is

\[ r \equiv \frac{n_B}{n_B'} = \frac{\kappa}{\kappa'} \cdot \frac{\sum_{k \neq 1} \frac{1}{M_k} \text{Im}\{3[(\lambda^\dagger \lambda)_{k1}]^2 + 2[(\lambda^\dagger \lambda)_{k1}(\xi^\dagger \xi)_{k1}]\}}{\sum_{k \neq 1} \frac{1}{M_k} \text{Im}\{3[(\xi^\dagger \xi)_{k1}]^2 + 2[(\xi^\dagger \xi)_{k1}(\xi^\dagger \xi)_{k1}]\}} , \] (6)

which means the baryon to dark matter relic abundance satisfies

\[ \frac{\Omega_{\text{baryon}}}{\Omega_{\text{dark matter}}} = \frac{m_N}{m_{DM}} \cdot r , \] (7)

where \( m_N, m_{DM} \) are the nucleon and dark baryon masses, respectively. For the resonant leptogenesis case, the factor \( \frac{1}{M_R} \) is replaced by \( \frac{1}{\Gamma_N} \).

Mirror symmetry implies that \( \epsilon = \epsilon' \) and \( \kappa = \kappa' \). This means \( r = 1 \), the ratio of relic abundance is totally determined by the mass ratio of nucleon and dark matter particle, i.e., \( m_N/m_{DM} \sim 1/5 \). We will see in the next section how the mass differences between the nucleons in the two sectors can arise dynamically.

There are several constraints on the parameters of the model so that an adequate lepton number is generated and the washout effects due to scattering and decays do not reduce these values below what is required by observations. We summarize them below (as noted, we have denoted the leptonic Yukawa couplings by \( h = \xi = \lambda \)).

- We need to make sure that the resulting lepton asymmetries do not get erased by scattering processes of the type \( \ell + H \leftrightarrow \bar{\ell} + \bar{H} \) via right-handed neutrino exchange \[10\]. These processes go like \( \frac{(h^\dagger h)^2 T^3}{M_R^2} \) and for them to be out of equilibrium at \( T \sim M_R \), we require \( M_R \geq \frac{(h^\dagger h)^2 M_{Pl}}{120\pi} \);

- The decays of the RH neutrinos contribute to the washout factor \( K \equiv \frac{\Gamma_N}{H(M_R)} \sim \frac{(h^\dagger h)M_{Pl}}{12\pi\sqrt{3}g_{*}M_R} \) and should be smaller than about \( 10^6 \);

- The amount of primordial lepton asymmetry for the hierarchical RH neutrino case is roughly given by \( \epsilon < \frac{h^\dagger h}{16\pi} \);

- Finally, we borrow a result from a subsequent section about mirror masses that in order to retain the success of BBN in our model, the mirror neutrinos must decay before the epoch of nucleosynthesis to ordinary leptons, which requires that we must have \( m_{\nu'} \geq 100 \text{ MeV} \). This translates into a constraint on the Yukawa couplings and masses of right-handed neutrinos as follows: \( \frac{h^2}{M_R} \geq \frac{0.1 \text{GeV}}{v_{w_k}} \).

The consequence of these constraints is that for the hierarchical RH neutrino case, it is not possible to get enough primordial lepton asymmetry while simultaneously satisfying the washout and mirror neutrino mass constraints. We therefore adopt the resonant leptogenesis with \( M_1 \approx M_2 \sim 10^8 \text{ GeV} \) in which case there is an enhancement factor of \( M_R/\Gamma_N \), so that it compensates for small Yukawa coupling effects.


B. The complete model

In this subsection, we describe the new features that go beyond the usual mirror models (i.e., SM and its duplicate in each sector) that are needed for consistency. In all the new features, we maintain the exact mirror symmetry for all dimension four terms. For simplicity we may add soft mirror symmetry breaking terms, which may arise from a mirror symmetric model at high scale via spontaneous symmetry breaking. The presence of the soft breaking terms allows us to have symmetry breaking patterns in the two sectors different while the interactions remain symmetric. The model presented is non-supersymmetric but its supersymmetric extension will preserve its main features.

Instead of one Higgs doublet as in the standard model, we consider two Higgs doublets \( H_u,d \) and one \( Y = 2 \) triplet \( \Delta \) to both sectors. And to avoid large flavor changing neutral current, a \( Z_2 \) symmetry is imposed in each sector so that the up-type fermions only couple to \( H_u \) or \( H'_u \) whereas the down-type fermions to \( H_d \) or \( H'_d \).

The triplet Higgs bosons couple to left-handed lepton pairs as

\[
L_T = -Y_\Delta \bar{\ell} e \Delta \ell - Y'_\Delta \bar{\ell}' e' \Delta' \ell' + h.c. \tag{8}
\]

The neutral components of the triplet Higgs bosons get vacuum expectation values (VEVs) after spontaneous breaking of electroweak symmetries. Therefore, the left-handed neutrinos in the two sectors obtain Majorana mass terms and the generic neutrino mass matrix is given as in eq. (13). We will assume that the triplet VEV (or type II) contribution to the mass of the known neutrinos is zero whereas it is non-zero for the mirror neutrinos.

A few comments about the Higgs sector and the vev of Higgs fields is in order. First, in order to get the mirror doublet vevs to be larger than the familiar SM Higgs, we need to include soft breaking mass terms e.g.

\[
L_{soft} = \sum_a M_a^2 H_a^\dagger H_a + \sum_a M'_a^2 H'_a^\dagger H'_a + m_T^2 \vec{T}^\dagger \cdot \vec{T} + m'_T^2 \vec{T'}^\dagger \cdot \vec{T'} \tag{9}
\]

where \( M_a^2, M'_a^2, m_T^2 \) are alls negative and \( |M'_a|^2 >> |M_a|^2 \). As a result, the mirror Higgs vevs become different from the familiar SM Higgs and its partner.

Note also that since the Higgs triplet of the visible sector has no vev, it is free of constraints from electroweak radiative correction constraints. Admittedly, such differences in vev require fine tuning of parameters, whose proper understanding is beyond the scope of this paper, where we discuss the new scenario for dark matter and not the naturalness of the values of all the parameters of the theory. In deriving our conclusions, we have taken into account all the allowed renormalizable couplings involving the Higgs fields. We don’t display them here for brevity since they do not affect our discussion in this work.

C. The mirror nucleons

Generally, the mass of the nucleon is composed by two parts, namely the trace anomaly part and the quark masses. The trace anomaly part is proportional to the hadronic scale \( \Lambda_{QCD} \) and the \( \beta \)-function of QCD, and the contribution from the quark masses are proportional to the quark masses. Take the proton as an example; we have

\[
m_p = k \Lambda_{QCD} + 2 h_{pu} m_u + h_{pd} m_d , \tag{10}
\]
where $h_{pu}$ and $h_{pd}$ can be determined from the pion-nucleon $\sigma$-term. In the SM, since $m_u$ and $m_d$ are much smaller than the hadronic scale, their contribution can be neglected. In the mirror sector the contribution of the mirror quark masses may not be negligible, so masses of the mirror proton and neutron can be written as

$$m_{p'} = k'_p \Lambda'_{QCD} + 2h'_{pu}m_u + h'_{pd}m_d,$$

$$m_{n'} = k'_n \Lambda'_{QCD} + 2h'_{nd}m_d + h'_{nu}m_u,$$

where we set different coefficients of proportionality for the mirror proton and neutron since if the masses of the mirror quarks are comparable with the mirror hadronic scale, the isospin symmetry is strongly broken and there is no reason to set them to be the same. Under exact mirror symmetry, to get the correct dark matter relic density one must have the mass relation: $m_{p'} \approx 5m_p$ for mirror proton as the dark matter particle, or $m_{n'} \approx 5m_p$ for dark neutron as dark matter. This difference in the nucleon masses in the two sectors can arise if the two electroweak scales are different with $v'_{wk} \gg v_{wk}$. What this difference of EW scales does is to make the mirror quarks ($t', b', c', \ldots$ etc.) much heavier than the familiar quarks of SM. If we further assume the two strong interaction coupling constants become equal at high scale due to mirror symmetry, we can get a relation between the two electroweak scales and the two hadronic scales using one loop evolution of the QCD couplings in the two sectors and assuming $\Lambda'_{QCD}$ and $\Lambda_{QCD}$ to be the scales where the QCD couplings $\alpha_{QCD}$ become of order one. Under different conditions for the mirror quark spectra, we get the following relations for the ratio $v_{wk}/v'_{wk}$.

$$\left(\frac{v_{wk}}{v'_{wk}}\right)^4 = \frac{\Lambda'^9_{QCD}(m_u m_d m_s)^{2/3}}{\Lambda'_{QCD}^9} \left(\frac{\sin \beta'}{\sin \beta}\right)^4 \left(\frac{\tan \beta'}{\tan \beta}\right)^{-2}, \quad \text{for } \Lambda'_{QCD} < m_{u'}, m_{d'};$$

$$\left(\frac{v_{wk}}{v'_{wk}}\right)^5 = \frac{\Lambda'^{27/2}_{QCD}(m_d m_s)}{\Lambda'_{QCD}^{27/2}} \left(\frac{\sin \beta'}{\sin \beta}\right)^5 \left(\frac{\tan \beta'}{\tan \beta}\right)^{-3}, \quad \text{for } m_{u'} < \Lambda'_{QCD} < m_{d'};$$

$$\left(\frac{v_{wk}}{v'_{wk}}\right)^5 = \frac{\Lambda'^{27/2}_{QCD}(m_u m_s)}{\Lambda'_{QCD}^{27/2}} \left(\frac{\sin \beta'}{\sin \beta}\right)^5 \left(\frac{\tan \beta'}{\tan \beta}\right)^{-2}, \quad \text{for } m_{d'} < \Lambda'_{QCD} < m_u;$$

$$\left(\frac{v_{wk}}{v'_{wk}}\right)^4 = \frac{\Lambda'^{27/2}_{QCD} m_s}{\Lambda'_{QCD}^{27/2}} \left(\frac{\sin \beta'}{\sin \beta}\right)^4 \left(\frac{\tan \beta'}{\tan \beta}\right)^{-2}, \quad \text{for } m_{u'}, m_{d'} < \Lambda'_{QCD} < m_{s'}.$$

These results do not depend on whether $m_{c'} < m_{u'}$ or $m_{c'} > m_{u'}$. From Eq. 12 one can see that $\Lambda'_{QCD}$ grows slowly with the increasing of $v'_{wk}$, that $\Lambda'/\Lambda = (v'_{wk}/v_{wk})^{1/3}/10$, there $v'$ increases a lot if $\Lambda'_{QCD}$ is a few times larger than $\Lambda_{QCD}$ which means $u'$ and $d'$ in the mirror nucleon might be nonrelativistic. In QCD, if the masses of $u$ and $d$ quarks were much larger than the hadronic scale they would be nonrelativistic inside the baryons. In that case the mass of the nucleon would be approximately equal to the sum of the quark masses plus a negative potential generated by the gluon field. Therefore, we can conclude that as $v'_{wk}$ grows larger and larger $h'_{u,d}$ approach 1 and $k'$ gets smaller and smaller. Using lattice QCD one can calculate $k'$ and $h'_{u,d}$ with different values of $v'_{wk}$. In this work, since we will see that the masses of $u'$ and $d'$ are comparable with the mirror hadronic scale, we assume that $h'_{pu} = h'_{pd} = h'_{nu} = h'_{nd} = 1$. For the other contribution, from the above analysis we know that as $v'_{wk}$ goes up it will get smaller, then crossover zero and then get negative, so in this work we will neglect this contribution.
In the SM, the neutron is slightly heavier than the proton due to $m_u < m_d$ and a free neutron will decay to a proton through beta decay. With the exact mirror Yukawa couplings, the dark neutron is expected to be heavier than the dark proton as in the SM. The situation could be different for two Higgs doublets $H^{(0)}_{u,d}$, with $\tan \beta = v_u/v_d$, $\tan \beta' = v'_u/v'_d$ different. $H^{(0)}_u$ only couples to up-type (mirror) quarks and (mirror) neutrinos while $H^{(0)}_d$ to down-type (mirror) quarks and charged (mirror) leptons, respectively, by imposed $Z_2$ symmetries in both the two sectors. As we can see below (Table I), when $\tan \beta'/\tan \beta > m_d/m_u$, the dark neutron is lighter than the dark proton.

In Table I, $\Lambda'_{QCD}$, the $u', d', s'$ and nucleon masses for different choices of $\tan \beta'$ are listed, while demanding the mass of the lightest nucleon in the mirror sector to be 5 GeV. It turns out $\Lambda'_{QCD}$ depends on $\tan \beta'$ very mildly, while the dark weak scale $v'_{wk}$ increases with $\tan \beta$. For low $\tan \beta'$, $m_{u'} < m_{d'}$, the dark proton is dark matter, and for larger $\tan \beta' \gtrsim 2 \tan \beta$, the dark neutron is dark matter. In the following calculation we will take the case $\tan \beta = 50$, $\tan \beta' = 200$. Therefore, the mirror neutron is the dark matter particle.

| $\tan \beta$ | $v'_{wk}$ (TeV) | $\Lambda'_{QCD}$ (GeV) | $m_{u'}$(GeV) | $m_{d'}$(GeV) | $m_{e'}$(GeV) | $m_{\mu'}$(GeV) | $m_{\nu'}$(GeV) |
|-------------|---------------|-----------------------|-------------|-------------|-------------|-------------|-------------|
| 50          | 123           | 1.06                  | 1.25        | 2.5         | 0.25        | 5.0         | 6.25        |
| 100         | 164           | 1.04                  | 1.67        | 1.67        | 0.17        | 5.0         | 5.0         |
| 200         | 246           | 1.06                  | 2.5         | 1.25        | 0.13        | 6.25        | 5.0         |
| 500         | 369           | 1.27                  | 3.57        | 0.71        | 0.07        | 7.85        | 5.0         |

TABLE I: The values of $v'_{wk}, \Lambda'_{QCD}$ for different $\tan \beta'$ with the light mirror nucleon mass fixed at 5 GeV. The other inputs are taken as $\tan \beta = 50$, $\Lambda_{QCD} = 200$ MeV, $m_u = 2.5$ MeV, $m_d = 5$ MeV.

III. NEUTRON MASSES

The neutrino mass matrix in the basis $(\nu, \bar{N'}, \nu')$ can be written as

$$
\mathcal{M} = \begin{pmatrix}
\mu & M_D & 0 \\
M_D^T & M_R & M_D^T \\
0 & M_D^T & \mu'
\end{pmatrix},
$$

where $M_D = \lambda v_{wk}$ and $M'_D = \xi v'_{wk}$, $\mu$ and $\mu'$ are Majorana mass matrices for SM neutrinos and mirror neutrinos generated from type-II seesaw mechanism, respectively. We take the exact mirror symmetry for Yukawa couplings, i.e., $\lambda = \xi = \hat{h}$. Since we want to roughly reproduce the general features of the neutrino masses and the mixing of the SM neutrinos, the flavor indices are suppressed in the following calculations.

From Table I $v'_{wk}$ is chosen about $10^3$ times larger than $v_{wk}$, so $M'_D$ is $10^3$ times larger than $M_D$ due to the mirror symmetry of the Yukawa couplings. If $\mu' M_R \ll M_D^2$, the SM neutrino mass lies in the type II + inverse seesaw regime, $M_\nu \approx \mu + M_D M_D^{-1} \mu' M_D M_D^{-1} M_D^T = \mu + \mu' (v_{wk}/v'_{wk})^2$, while the mass of the dark sector neutrino receives dominant contribution from type I seesaw mechanism, i.e., $M'_\nu \approx \mu' - M'_D M_R^{-1} M_D^T \approx -M'_D M_R^{-1} M_D^T$. On the other hand, if $\mu' M_R > M_D^2$, the SM neutrino is contributed from type I and type II seesaw mass $M'_\nu \approx \mu' - M_D M_R^{-1} M_D^T$, while the dark neutrino mass is mainly type II, $M'_\nu \approx \mu'$. However, in the latter case, the triplet Higgs contribution dominates the leptogenesis in
the dark sector \[12\] which will ruin the relation \(n_B = n_{B'}\). Hence, in this work we adopt the first scenario. Also, since \(\mu\) is determined by the interaction between doublet Higgs bosons and the triplet Higgs boson in the SM sector, it is independent of other quantities in the neutrino mass matrix; therefore, we can simply assume it to be much smaller than the neutrino masses and neglect its contribution.

In order to generate the relic density of baryon and dark matter, and to avoid the constraints from various experiments and observations which will be discussed in the following sections, we choose the following parameters:

\[
M_R \approx 1 \times 10^8 \text{GeV}, \quad h = 0.015, \quad v_{wk}'/v_{wk} = 10^3, \quad \mu' \approx 100\text{KeV}.
\] (14)

The resulting neutrino masses are of order \(M_\nu \approx 0.1\text{ eV}, \ M'_\nu \approx 150\text{ MeV}\). We note that there is a non-trivial mixing between the SM and dark neutrinos which is \(\hat{\nu} = \nu + U_{\nu\nu'}\nu'\), \(U_{\nu\nu'} \approx M_D/M'_D \approx 1.0 \times 10^{-3}\); i.e., the mixing between familiar and mirror neutrinos is a universal number for all flavors in the flavor basis. The dark neutrino, once produced, can decay to \(e^+e^-\nu\) due to this mixing through weak interactions. For the parameters set in Eq. (14) the lifetime of dark neutrinos can be estimated as \(\tau_{\nu'} < 0.5\text{ sec}\). Therefore the mirror neutrinos decay to the SM sector before the BBN epoch.

We wish to emphasize that the set of parameters in Eq. (14) is quite unique. When we fix the ratio of \(v_{wk}/v_{wk}'\) from Eq. (12), the SM neutrino masses and mixings are completely determined by \(\mu'\). Furthermore, we cannot make \(M_R\) heavier because that will reduce the mass for dark neutrino thereby contradicting with the constraints from BBN. On the other hand, for \(M_R = 10^8\text{ GeV}\), we have the \(K\) factor relevant for leptogenesis \(K = \frac{M_D M_R^{-1} M_T}{m_*} \approx 1.3 \cdot 10^{-7}\text{GeV/10}^{-3}\text{eV} = 1.3 \times 10^5\), which implies a washout factor \(\kappa \approx 1 \times 10^{-6}\). According to Eqs. (4) and (5), to get \(\eta_b \approx 5 \cdot 10^{-10}\), we need to have a value of primordial CP asymmetry \(\epsilon \sim 0.05\). In this scenario, as already noted, leptogenesis can be realized if at least two right-handed neutrinos are quasi-degenerate \([7]\). Decreasing \(M_R\) will increase the \(K\) factor; larger \(K\) implies stronger washout effect or smaller \(\kappa\), threatening the success of leptogenesis.

IV. THE CONSTRAINTS

A. Constraints from BBN

BBN is well constrained by the expansion rate of the universe at the temperature \(T_{BBN} \sim 1\text{ MeV}\). The Hubble expansion rate is determined by the total energy density, which constrains the light degrees. When at \(T \sim 1\text{ MeV}\), we have the number of degrees of freedom (d.o.f.) \(g_* = 10.75\) contributed from the SM photons, electrons and neutrinos. The constraints on new d.o.f. is conventionally quoted as \(\Delta N_\nu\), the effective number of additional light neutrino species. A reliable bound is \(\Delta N_\nu \leq 1.44\) at 95\% CL by various present observations \([13]\).

In the symmetric mirror model, the mirror neutrinos as well as the mirror photon and electrons will contribute another 10.75 to \(g_*\), thereby completely spoiling the success of BBN. One way to avoid the BBN bound is to have the hidden sector with lower temperature than the SM sector. This could be achieved if the reheating temperatures after inflation the two sectors are different\([14]\). However, this scenario does not work here since the SM sector and the mirror sector are connected by the common right-handed neutrinos which can bring
the two sectors back into thermal equilibrium even if they have different couplings with
the inflaton. When the temperature is lower than the mass of the lightest right-handed
neutrino, the interaction between the Higgs bosons in two sectors $\lambda |H|^2 |H'|^2$ and the kinetic
mixing between the $U(1)_Y$ gauge bosons $\epsilon^{\mu\nu} B^\mu B'^\nu$ (see Appendix) provide two alternative
mechanisms to keep the two sectors in thermal equilibrium. The kinetic mixing of the $U(1)_Y$
bosons of the two sectors $B$ and $B'$ induces the kinetic mixing between the photon and mirror
photon; therefore if mirror photon is massless there is a long range interaction between the
two sectors. As a result the two sectors will never be thermally decoupled from each other
thereby increasing the number of degrees of freedom at BBN to unacceptable values.

In our model, since the mirror neutrinos have masses above a 100 MeV and lifetimes
much less than a second, they do not pose any problem for BBN. We only have to make
sure that the photon contribution is eliminated. In order to achieve this, we work with two-
Higgs-doublets in each sector such that the $\mu$Higgs-doublet in each sector is, of course, kept unbroken. We
choose the mass of the mirror photon to be $O(100)$ MeV.

The massive dark photon is coupled to the SM fermions through the kinetic mixing $\epsilon_{\gamma}^2 F^{\mu\nu} F'^{\mu
\nu}$ and therefore decays to $e^\pm$ pair. The decay rate is given by $\Gamma_{\gamma'} = \alpha_{em} \epsilon_{\gamma}^2 m_{\gamma'}/3$.
The lifetime of the dark photon is

$$\tau_{\gamma'} \approx \left(\frac{50\text{MeV}}{m_{\gamma'}}\right) \left(\frac{7 \times 10^{-11}}{\epsilon_{\gamma}}\right)^2 \text{sec}. \quad (15)$$

For $m_{\gamma'} = 50 \text{ MeV}$, $\epsilon_{\gamma} > 7 \times 10^{-11}$ is needed to make the dark photon lifetime be shorter
than 1 sec.

QED precision measurements provides constraints on the coupling $\epsilon_{\gamma}$. The most important
constraint comes from the measurement of the muon magnetic moment, which gives an
upper bound $\epsilon_{\gamma}^2 < 2 \times 10^{-5}(m_{\gamma'}/100\text{MeV})^2$ [16].

The mass of the mirror photon is induced by the nonvanishing VEVs of charged compo-
nents of the mirror doublet Higgs bosons. The other non-trivial consequence of breaking
dark $U(1)'_{em}$ is the charged mirror particles will mix with neutral particles since $\langle H'^\pm \rangle \neq 0$.
These small mixings allow the mirror proton to radiatively decay to mirror neutron, i.e., $p' \rightarrow n'\gamma'$ and the mirror electron to decay as $e' \rightarrow \gamma' \nu$. The size of this mixing can roughly be estimated as $U_{p'n'}^\dagger \sim \frac{\nu_{e_{\gamma}Y_d}^\dagger}{\nu_{w_{\gamma}Y_N}} \approx \frac{100\text{MeV}y_{\gamma}}{100\text{MeV}y_u} \approx 10^{-6} \cdot 2 \tan \beta$ for the case we are interested. For
$\tan \beta \sim 50$, this gives $U_{p'n'} \sim 10^{-4}$.

The decay rate for the process $p' \rightarrow n'\gamma'$ is given by

$$\Gamma \approx \frac{\alpha_{em} U_{p'n'}^2 m_{p'}}{4} F\left(\frac{m_{n'}}{m_{p'}}, \frac{m_{\gamma'}}{m_{p'}}\right), \quad (16)$$

where $F(x, y) = \frac{(1-x^2+y^2)(1-x^2-y^2)}{y^2} \left[(1 - (x + y)^2)(1 - (x - y)^2)\right]^{1/2}$. For the case $\tan \beta' = 200$, one can obtain the lifetime of mirror proton $\tau_{p'} \approx 10^{-15}$ sec. For the mirror electron decay, $e' \rightarrow \gamma' \nu$, inputing $m_{\nu} = 0.13 \text{ GeV}$ and $U_{e'\nu} \approx \frac{U_{e'\nu}}{U_{\nu_{YD}}} \approx 10^{-9}$ which agrees with
the numerical result, its lifetime is estimated to be $5 \times 10^{-2}$ sec. Both of them decay before
the epoch of BBN.
B. Constraints from neutrinoless double beta decay

The dark neutrino contributes to the $0\nu\beta\beta$ decay process by exchange $\nu'$ due to the mixing with the active neutrino $\nu_e$. Conventionally one parameterizes the experimental bounds on $0\nu\beta\beta$ decay as a limit on the “effective” neutrino mass $m_{\nu'}^{\text{eff}}$. For light neutrinos, $m_{\nu'}^{\text{eff}}$ is defined as $\Sigma_i U_{\nu_i}^2 m_i$. The current upper bound is $m_{\nu'}^{\text{eff}} < 0.5$ eV from the $^{76}\text{Ge}$ $0\nu\beta\beta$ experiment [17].

When the dark sterile neutrino masses are heavy compared to $\mathcal{O}(100)$ MeV, like the case we are considering here, one gets the contribution to the effective mass [18]

$$m_{\nu'}^{\text{eff}} = \frac{U_{\nu\nu'}^2 q_F^2}{3m_{\nu'}} ,$$

(17)

where $q_F$ is the nucleon Fermi momentum and its value is taken as $q \approx 60$ MeV [18]. So, inputing $U_{\nu\nu'} = 10^{-3}$, $m_{\nu'} = 136$ MeV, one finds that this parameter choice has tension to satisfy the experimental limit. However, if there is cancellation between the contributions from $\nu'$'s and light neutrinos, the constraint could be relaxed. Furthermore, considering the structure of the total neutrino mass matrix in Eq. (13), we notice that the light SM neutrino mass matrix is $M_\nu = \mu + \mu'(v_{\text{wk}}/v_{\text{wk}})^2$ due to the exact mirror symmetric Yukawa couplings. This means there is no family-mixing between the light neutrinos and the sterile dark neutrinos, $U_{ij'} \propto \delta_{ij'}$ as noted already. The mirror neutrino contribution to $0\nu\beta\beta$ decay is proportional to the $e'e'$ component of the sterile dark neutrino mass matrix $(M'_D M^{-1}_R M'_D^T)_{e'e'}$. The active neutrino mass matrix, on the other hand, is given by the matrix $\mu'$. Therefore they are unrelated and we are free to choose the $e'e'$ component of the mirror neutrino mass matrix without affecting the active neutrino mixings. Choosing a tiny or even vanishing value for this element can guarantee the model to avoid the experimental limits.

C. Constraints of self interaction cross-section of dark matter

In our model since mirror neutron is the dark matter, it will have strong scattering against other mirror neutrons. There are two sources for this scattering to arise from: (i) strong scattering and (ii) electromagnetic scattering. As regards strong scattering, we note that familiar low energy neutron scattering of neutrons off protons has a cross-section of order $\sigma_{np} \sim 10^{-24}$ cm$^2$ and by isospin symmetry, the $\sigma_{nn} \simeq \sigma_{np}$. Since in the mirror sector, $\Lambda'_{QCD} \sim 10\Lambda_{QCD}$, we expect the cross section $\sigma_{n'n'} \sim 10^{-26}$ cm$^2$. Note that the upper bound on the dark matter self interaction cross section is given by [19] $\sigma/m_{\text{dark matter}} < 1.25$ cm$^2$ gram$^{-1}$ which can interpreted as $\sigma_{n'n'} \leq 10^{-23}$ cm$^2$ in our case where the dark matter mass is about 5 GeV. Coming to the electromagnetic scattering cross-section, we note that $n'$ has no mirror electric charge but a mirror magnetic moment given roughly by $\mu_{n'} \sim e k/m_{n'}$, where $k$ can be estimated as $m_{\text{dark matter}} \times v$, and the velocity $v/c$ is roughly $10^{-3}$. Therefore the electromagnetic vertex which goes into self scattering has a strength of order $10^{-3}e$. Now, since the momentum transfer is much smaller than the mass of the dark photon, the vertex of the four-dark matter interaction can be written as $10^{-6}e^2/m_{\gamma'}^2$. The cross section is estimated to be $(10^{-6}e^2/m_{\gamma'}^2) \times m_{n'}^2/(4\pi)$ which is at most $10^{-35}$ cm$^2$, which is much below the bullet cluster upper bound.
V. DIRECT DETECTION

Asymmetric dark matter can be detected directly by observing the nucleus recoil at low background experiments. The effective operators for DM-nucleon interaction can be generalized as

\[ \chi \Gamma_1 \chi N \Gamma_2 N \],

where \( \Gamma_1, \Gamma_2 \) can be the combinations of scalar, pseudoscalar, vector, axial, tensor or pseudo-tensor [20]. In the non-relativistic limit, all the interactions can be reduced to two terms, spin-independent (SI) and spin-dependent (SD).

In our model, the direct detection process can be induced by the interaction between the Higgs bosons in the two sectors and kinetic mixing between the gauge bosons.

First, let’s consider the direct detection via the Higgs interaction \( f|H|^2|H'|^2 \). After the spontaneous symmetry breaking in the two sectors, this term generates an effective four-fermion interaction between nucleons and mirror nucleons, which can be written as

\[ L_{\text{eff}} = \delta f n' n' NN \],

(19)

where \( \delta f \approx \frac{(f_m N m_n')}{(10 m_h^2 m_h')} \) [21]. The total cross section of the elastic scattering between \( n' \) and the nucleon can be written as

\[ \sigma_{n'n} \approx \frac{\mu_r^2}{\pi} \delta f^2 \approx 10^{-29} f^2 \text{ GeV}^{-2} \approx 10^{-57} f^2 \text{ cm}^2 \],

(20)

where \( \mu_r = m_N m_{n'}/(m_N + m_{n'}) \) is the reduced mass of \( n' \) and nucleon \( N \), \( m_h \) is the mass of a Higgs boson from the SM sector which is set to be 100 GeV and \( m_{h'} \) from the mirror sector is set to be 100 TeV. This is well below the present upper bound of \( 10^{-39} \text{ cm}^2 \) for around 5 GeV dark matter.

The kinetic mixing between photon and mirror photon plays an important role in the direct detection. The local velocity of dark matter is assumed to be \( \sim 200 \text{ km/s} \), so for a 5 GeV dark matter its kinetic energy is about 1 KeV, which is much smaller than the mass of the mirror photon which is assumed to be 50 MeV. Therefore, the interaction between the nucleon and the mirror neutron can also be viewed as a point-like interaction. The mirror neutron interacts with the mirror photon through its anomalous magnetic moment and also through its mixing with the mirror proton due to the breaking of the mirror electromagnetic \( U(1) \) symmetry.

Since \( u' \) and \( d' \) are heavy, one can use constituent quark model to estimate the magnetic moment of the mirror neutron.

\[ \mu_n' = \frac{4}{3} \mu_d' - \frac{1}{3} \mu_u' \],

(21)

where \( \mu'_{(u,d)} = eQ_{u,d}/2m_{u,d} \) are magnetic moments of \( u' \) and \( d' \) where \( Q_u = 2/3 \) and \( Q_d = -1/3 \) are the mirror charges of \( u' \) and \( d' \). In the nonrelativistic regime the total cross section between the nucleon and the mirror neutron induced by this magnetic interaction can be estimated as

\[ \sigma_{n'n}^{\text{mag}} \approx \frac{\alpha_{em}^2 \pi \varepsilon_\gamma^2 |\vec{q}|^4}{(m_{n'} + m_N)^2 m_{n'}^4} \lesssim 4 \times 10^{-8} \varepsilon_\gamma^2 \text{ GeV}^{-2} \lesssim 10^{-35} \varepsilon_\gamma^2 \text{ cm}^2 \],

(22)

where \( \vec{q} \) is the momentum transfer during the collision which is about 5 MeV. A less important contribution comes from the mixing between the \( n' \) and \( p' \). The cross section induced
by this mixing can be written as

\[ \sigma_{\gamma'\gamma'} \approx \frac{\mu^2}{\pi} \left( \frac{\varepsilon_\gamma \varepsilon_\gamma' U_{\gamma'\gamma}}{m_{\gamma'}^2} \right)^2 \approx 10^{-14} \varepsilon_\gamma^2 \text{ GeV}^{-2} \approx 10^{-41} \varepsilon_\gamma^2 \text{ cm}^2. \]  

Recalling the upper bound on \( \varepsilon_\gamma \) from QED precision measurements given above, we conclude that the cross section is well below the upper bound set by direct detection experiments and our dark matter could be accessible to direct search experiments in future.

From Eq. (A2) in Appendix, one can see that interactions through \( Z \) and \( Z' \) between the two sectors are either suppressed by the nature of the magnetic interaction or by the \( M_Z^2/M_{Z'}^2 \), and then are negligible.

**VI. CONCLUSIONS**

In this work, we have proposed that the dark matter of the Universe be identified with the lightest baryon of a possible mirror duplicate of the standard model with the only difference between the two sectors being in the symmetry breaking patterns. Prior to spontaneous symmetry breaking, this model has no free parameters due to mirror symmetry. The lightest dark nucleon is stable due to the mirror analog of baryon number and becomes the dark matter. It is an asymmetric dark matter with its anti-mirror baryon part suppressed in a manner analogous to the matter-anti-matter asymmetry in the standard model sector. The introduction of a common set of right-handed neutrinos connecting the two sectors allows a common mechanism for the genesis of the matter-anti-matter asymmetry in both sectors thereby helping us to understand why the dark matter and normal baryon contribution to the energy density of the Universe are not too different from each other. One only has to make the assumption that the dark nucleon is five times heavier than the familiar nucleon – an assumption that is easily understood if the electroweak scales in the two sectors are different. We show that this model can be consistent with the constraints of BBN and neutrinoless double beta decay. The mirror photon in our model is massive but mixes with the normal photon to avoid the BBN constraints.

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**Appendix A: Kinetic mixing between gauge bosons**

The kinetic mixing between the \( U(1) \) gauge bosons \( B \) and \( B' \) will induce kinetic mixings of \( \gamma \) and \( Z \) with their mirror partners due to the spontaneous symmetry breaking in both sectors. Therefore the Lagrangian for kinetic mixing can be written as

\[ \mathcal{L}_{\text{kin}}^{\text{mix}} = \frac{\varepsilon_B}{2} \left[ \cos^2 \theta_W F^{\mu\nu} F'_{\mu\nu} - \sin \theta_W \cos \theta_W (F^{\mu\nu} Z'_{\mu\nu} + Z^{\mu\nu} F'_{\mu\nu}) + \sin^2 \theta_W Z^{\mu\nu} Z'_{\mu\nu} \right], \]  

(A1)
where $\theta_W$ is the Weinberg angle. Since exact mirror symmetry is assumed the two sectors share the same Weinberg angle at tree level. Using the fact that $M_{Z'}^2 \gg M_Z^2 \gg m_\gamma^2$, to the leading order of $\epsilon_B$ the following redefinitions of $A_\mu$, $A'_\mu$, $Z_\mu$ and $Z'_\mu$ diagonalize both the kinetic terms and the mass terms of gauge bosons:

\begin{align}
A_\mu & \longrightarrow A_\mu + \epsilon_\gamma A'_\mu - \sqrt{\epsilon_\gamma \epsilon_Z} Z'_\mu ; \\
Z_\mu & \longrightarrow Z_\mu + \epsilon Z Z'_\mu + \sqrt{\epsilon_\gamma \epsilon_Z} \frac{M_{Z'}^2}{M_{Z'}^2} A'_\mu ; \\
A'_\mu & \longrightarrow A'_\mu - \sqrt{\epsilon_\gamma \epsilon_Z} Z_\mu ; \\
Z'_\mu & \longrightarrow Z'_\mu - \epsilon Z \frac{M_Z^2}{M_{Z'}^2} Z_\mu ,
\end{align}

(A2)

where $\epsilon_\gamma = \epsilon_B \cos^2 \theta_W$ and $\epsilon_Z = \epsilon_B \sin^2 \theta_W$.

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