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Originally Asymmetric Dark Matter

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Abstract

We propose a scenario with a fermion dark matter, where the dark matter particle used to be the Dirac fermion, but it takes the form of the Majorana fermion at a late-time. The relic number density of the dark matter is determined by the dark matter asymmetry generated through the same mechanism as leptogenesis when the dark matter was the Dirac fermion. After efficient dark matter annihilation processes have frozen out, a phase transition of a scalar field takes place and generates Majorana mass terms to turn the dark matter particle into the Majorana fermion. In order to address this scenario in detail, we propose two simple models. The first one is based on the Standard Model (SM) gauge group and the dark matter originates the $SU(2)_L$ doublet Dirac fermion, analogous to the Higgsino-like neutralino in supersymmetric models. We estimate the spin-independent/dependent elastic scattering cross sections of this late-time Majorana dark matter with a proton and find the possibility to discover it by the direct and/or indirect dark matter search experiments in the near future. The second model is based on the $B-L$ gauged extension of the SM, where the dark matter is a SM singlet. Although this model is similar to the so-called Higgs portal dark matter scenario, the spin-independent elastic scattering cross section can be large enough to detect this dark matter in future experiments.

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I. INTRODUCTION

The origin of baryon asymmetry as well as of dark matter (DM) are fundamental questions in cosmology. Until now various mechanisms for the generation of baryon asymmetry and the candidates for dark matter have been investigated. Under the standard paradigm, the dark matter is made up with the weakly interacting massive particle (WIMP) which freezes out from the thermal equilibrium, whilst the baryon asymmetry is supposed to be generated by a mechanism which has nothing to do with the relic abundance of the WIMP dark matter. However, the observed baryon density Ω_b and that of dark matter Ω_{DM} are close to each other [1], namely $\Omega_{DM}/\Omega_b \sim 5$, and this fact might imply a certain connection between the two densities in history of the Universe. For example, the Q-ball generated in the early Universe [2–7] can be a single source to explain this concordance.

The so-called "asymmetric dark matter" (ADM) scenario offers another approach, where similarly to the relic baryon density, the dark matter relic density is determined by a nonvanishing asymmetry between dark matter particle and its anti-particle [8–10], rather than the usual thermal freeze-out via annihilation. In many ADM models proposed so far [11–18], the dark matter particle shares a common quantum number with the baryon and the dark matter asymmetry is generated associated with the baryogenesis. For a successful ADM scenario, the following two conditions should be satisfied. First, the dark matter particle is a complex field and "dark matter number" can be defined which distinguishes dark matter particle and its anti-particle. Second, the annihilation cross section between dark matter and anti-dark matter particles is large enough for the symmetric component of the abundance to be almost completely erased [19–23]. As a result, the relic abundance of dark matter is determined by the dark matter-anti-dark matter asymmetry left in the Universe. In some of ADM models, this condition of efficient annihilation processes is achieved in the presence of light mediator particles [18, 24].

If an ADM has the electroweak interaction, especially the coupling with Z-boson with a large annihilation cross section, the dark matter direct search experiments [25] give a very severe constraint on mass of the dark matter $\gtrsim \mathcal{O}(10\ \text{TeV})$. This constraint can be avoided if a tiny dark matter number violating mass term is introduced. However, it has been found that this term should be very small $\lesssim 10^{-41}$ GeV, in order to suppress the dark matter-anti-dark matter oscillation in the early Universe [26–29]. While such a tiny dark matter number violation may be attractive because the dark matter indirect detection becomes available [26–28], its smallness gives rise to the fine-tuning issue in theoretical point of view.

In this paper, we propose a novel ADM scenario, in which the dark matter relic density is determined by the dark matter asymmetry in the early Universe, but the nature of the particle is Majorana without introducing such the extremely tiny parameter as mentioned above. The key ingredient is that the dark matter is the Dirac fermion in the very early Universe, but a late-time phase transition generates Majorana mass terms and as a result, the initially Dirac fermion dark matter takes the form of the Majorana fermion, after the symmetric part of the DM abundance has been erased by the annihilation between dark matters and anti-dark matters. In this paper we propose two simple models to address this scenario in detail. In the first model (Model A) the dark matter particle has the electroweak interaction through which the dark matter and anti-dark matter particles annihilate very efficiently in the early Universe. We consider the gauged $U(1)_{B-L}$ extension of the Standard Model for the second model (Model B), where the dark matter particle is singlet under the SM gauge group but has a $U(1)_{B-L}$ charge. For completeness, we also consider a concrete mechanism for generating both the baryon and dark matter asymmetries. Various scenarios for baryogenesis by means of, for example, the out-of-equilibrium decay of heavy particles [30–33], a scalar condensate [34–37] and a CPT violating system [38, 39] have been proposed in the framework of the ADM scenario. Similarly to the scenario proposed in Ref. [33], we adopt the same mechanism as leptogenesis [40], which is probably the simplest scenario for baryogenesis, for generating the dark matter asymmetry as well as the baryon asymmetry in the Universe.

The paper is organized as follows. We first consider Model A in the next section, where new particles (Dirac fermions and complex scalars) of the electroweak doublets and singlets are introduced. In the early Universe, the out-of-equilibrium decay of the right-handed neutrinos generates the dark matter asymmetry as well as the baryon asymmetry. The Dirac fermion dark matters efficiently annihilate through the electroweak interaction and the resultant relic abundance is determined by the dark matter asymmetry. A late-time phase transition by a light scalar generates Majorana mass terms, by which the initially Dirac fermion dark matter has taken the form of the Majorana particle. We estimate the elastic scattering cross sections of the dark matter particle with a proton which are relevant for the direct/indirect dark matter search experiments. In Sec. III, we introduce Model B based on the gauged $U(1)_{B-L}$ extension of the SM. New particles introduced in this model are all the SM gauge singlets. Our discussion on Model B is analogous to that on Model A, but in Model B efficient annihilation processes of Dirac fermion dark matters are realized by the resonance effect of intermediate particles such as the $B-L$ gauge boson and Higgs bosons. Although this model is similar to the so-called Higgs portal dark matter scenario, we will see a crucial difference. Sec. IV is devoted for summary and discussions. Our notations and the formulas used in our analysis are listed in Appendix.

II. MODEL A

A. Setup

This model is based on the SM gauge group. We introduce a vector-like pair of $SU(2)_L$ doublet fermions, a vector-like pair of SM singlet fermions, and SM singlet and doublet complex scalars, as well as the right-handed neutrinos. The charge assignments are shown in Table I. Here, we have introduced a discrete Z_2 parity to guarantee the stability of dark matter particle and a global $U(1)$ symmetry to forbid Majorana mass terms for the gauge singlet fermions. The new doublet fermions are similar to the Higgsinos in the minimal supersymmetric Standard Model (MSSM) while the singlet fermions behaves like a pair of Bino without Majorana masses. The dark matter particle, which turns out to be a mixture of the doublet and singlet fermions after the electroweak symmetry breaking, has a large annihilation cross section through the electroweak gauge interaction.

The gauge invariant Lagrangian relevant to our discussion is given by [1]

$$\begin{aligned} -\mathcal{L} \supset\ & \frac{1}{2} y \bar{\chi_R^c} \phi \chi_R + \frac{1}{2} y \bar{\chi_L^c} \phi \chi_L + \frac{1}{2} \bar{N_R^c} M_N N_R \\ & + \mu_\chi \bar{\chi_L} \chi_R + \mu_\psi \bar{\psi_L} \psi_R + Y \bar{\psi_L} \Phi \chi_R + Y \bar{\chi_L} \Phi^\dagger \psi_R + Y_N \bar{\psi_L} \eta^\dagger N_R + h.c. \\ & + V(\Phi, \eta, \phi) \end{aligned} \tag{1}$$

with

$$\begin{aligned} V(\Phi, \eta, \phi) =\ & m_1^2 |\Phi|^2 + \lambda_1 |\Phi|^4 + m_3^2 |\phi|^2 + (\lambda_4 |\Phi|^2 + \lambda_6 |\eta|^2) |\phi|^2 + \lambda_7 |\phi|^4 \\ & + m_4^2 |\eta|^2 + \lambda_8 |\Phi|^2 |\eta|^2 + \lambda_{10} |\eta|^4 + h.c., \end{aligned} \tag{2}$$

[1] To make our discussion simple, we have dropped a term $Y'_N \bar{N_R^c} \eta \psi_R$, assuming $Y'_N \ll Y_N$.

TABLE I: Particle contents of Model A

Fields	$SU(2)_L$	$U(1)_Y$	Z_2-parity	Global U(1)
ψ_L	2	+1/2	−	−1
ψ_R	2	+1/2	−	−1
χ_L	1	0	−	−1
χ_R	1	0	−	−1
Φ	2	+1/2	+	0
η	2	−1/2	−	+1
ϕ	1	0	+	+2
N_R	1	0	+	0

where Φ denotes the SM Higgs doublet field. For simplicity, we have assumed common Yukawa coupling constants (y, Y and Y_N) between ϕ and $\chi_{L,R}$, between Φ, $\chi_{L,R}$ and $\psi_{R,L}$, and between η, $\psi_{L,R}$ and N_R, respectively. M_N is the Majorana mass term of the right-handed neutrinos N_R. We have omitted the flavor index for the right-handed neutrinos.

We adjust the parameters in the scalar potential so as for Φ and ϕ (but not η) to develop vacuum expectation values (VEVs) at a true vacuum, where we expand the Higgs fields as

$$\Phi = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}(v+h) \end{pmatrix}, \quad \phi = \frac{1}{\sqrt{2}}(v_s + \varphi), \tag{3}$$

with VEVs, $v = 246$ GeV and v_s. The mass eigenstates are obtained by the unitary rotation

$$\begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = \begin{pmatrix} S_{11} & S_{12} \\ S_{21} & S_{22} \end{pmatrix} \begin{pmatrix} \varphi \\ h \end{pmatrix} = \begin{pmatrix} \cos\alpha & \sin\alpha \\ -\sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} \varphi \\ h \end{pmatrix}, \tag{4}$$

with α being the mixing angle, which diagonalizes the mass matrix

$$\text{Mass term } = \frac{1}{2}(\varphi \ h) \begin{pmatrix} 2v_s^2\lambda_7 & vv_s\lambda_4 \\ vv_s\lambda_4 & 2v^2\lambda_1 \end{pmatrix} \begin{pmatrix} \varphi \\ h \end{pmatrix}. \tag{5}$$

In the following, we consider the phase transitions by VEVs of Φ and ϕ which take place in the early Universe as the temperature of the Universe is decreasing according to the expansion. Furthermore, we concentrate on a parameter space where the electroweak symmetry breaking occurs much earlier than the development of non-zero ϕ VEV.

B. Cosmological evolution

The right-handed neutrinos N_R have two decay modes; decay into lepton and Higgs ($N_R \to L\Phi$), and decay into η and ψ ($N_R \to \psi\eta$). Lepton asymmetry can be generated by the former decay mode as in usual thermal leptogenesis [40]. The resultant asymmetry is expressed as

$$\frac{n_L}{s} = \frac{\kappa_L}{g_*}\epsilon_{(N_R\to L\Phi)} \times \mathrm{Br}(N_R \to L\Phi) \tag{6}$$

by using the dilution factor κ_L corresponding to the washout effect for the generated lepton asymmetry, the branching ratio of N_R decay into lepton $\mathrm{Br}(N_R \to L\Phi)$ and the CP asymmetry parameter in the decay

$$\epsilon_{(N_R\to L\Phi)} = \frac{\Gamma(N_R \to L\Phi) - \Gamma(N_R \to \bar{L}\Phi^*)}{\Gamma(N_R \to L\Phi) + \Gamma(N_R \to \bar{L}\Phi^*)}. \tag{7}$$

The same mechanism generates the asymmetry between ψ and $\bar{\psi}$ such as

$$\frac{n_\psi}{s} = \frac{\kappa_\psi}{g_*}\epsilon_{(N_R\to\psi\eta)} \times \mathrm{Br}(N_R \to \psi\eta). \tag{8}$$

Here κ_ψ is the corresponding dilution factor, the branching ratio of N_R decay into ψ is $\mathrm{Br}(N_R \to \psi\eta)$, and the CP asymmetry in the decay is defined as

$$\epsilon_{(N_R\to\psi\eta)} = \frac{\Gamma(N_R \to \psi\eta) - \Gamma(N_R \to \bar{\psi}\eta^*)}{\Gamma(N_R \to \psi\eta) + \Gamma(N_R \to \bar{\psi}\eta^*)}. \tag{9}$$

In the following, the dark matter particle is the Dirac fermion which is a mixture of ψ and χ and therefore, this ψ-asymmetry generated by N_R decay is nothing but the asymmetry of ψ dark matter sector. Notice that the asymmetry between η and η^* is also generated.

After the electroweak symmetry breaking but ϕ has not developed its VEV yet, the mass matrix for the Z_2-odd Dirac fermions is given by

$$(\bar{\chi_L}\bar{\chi_R^c}\bar{\psi_L}\bar{\eta_R^c})M \begin{pmatrix} \chi_L^c \\ \chi_R \\ \psi_L^c \\ \psi_R \end{pmatrix}, \tag{10}$$

where

$$M = \begin{pmatrix} 0 & \mu_\chi & 0 & Y\frac{v}{\sqrt{2}} \\ \mu_\chi & 0 & Y\frac{v}{\sqrt{2}} & 0 \\ 0 & Y\frac{v}{\sqrt{2}} & 0 & \mu_\psi \\ Y\frac{v}{\sqrt{2}} & 0 & \mu_\psi & 0 \end{pmatrix}, \tag{11}$$

so that the mass eigenvalues are found to be

$$M_{\chi\psi}^{Dirac} = \frac{\mu_\chi + \mu_\psi \pm \sqrt{2Y^2v^2 + (\mu_\chi - \mu_\psi)^2}}{2}. \tag{12}$$

Provided η is heavier than those fermions, the lightest mass eigenstate among the Dirac fermions is the dark matter particle. Since this Dirac fermion partly consists of the $SU(2)$ doublet fermions (ψ), its annihilation processes are efficient in particular when the annihilation channels into WW and/or ZZ are kinematically allowed. This situation is analogous to the Higgsino-like neutralino dark matter in the MSSM. Therefore, Dirac fermion dark matters and anti-Dirac fermion dark matters can almost completely annihilate and only the dark matter asymmetry is left in the Universe. Similarly, the same is true for η and η^*, hence its symmetric component of η and η^* is erased by annihilation through $SU(2)$ gauge interactions.

Since the ϕ is light, the scalar potential at zero-temperature is simplified at low energies where the other scalars (and fermions) are decoupled:

$$V(\Phi, \eta, \phi) \ \sim\ -m_\phi^2|\phi|^2 + \lambda_7|\phi|^4. \tag{13}$$

Here we have fixed parameters $m_\phi^2 = -m_3^2 - \lambda_4 v^2/2 > 0$ and hence ϕ develops the VEV,

$$\langle\phi\rangle = \frac{v_s}{\sqrt{2}} = \sqrt{\frac{m_\phi^2}{2\lambda_7}}. \tag{14}$$

In the early Universe, this scalar potential is modified by the thermal effect and temperature corrections are given by [41]

$$\delta V = \frac{1}{64\pi^2}m^4\left(\ln\frac{m^2}{\mu^2} - \frac{3}{2}\right) + \frac{T^4}{2\pi^2}J_B(m^2/T^2), \tag{15}$$

where $m^2 = -m_\phi^2 + 4\lambda_7|\phi|^2$. In a high-temperature, the function $J_B(m^2/T^2)$ is expanded as

$$J_B(m^2/T^2) = -\frac{\pi^4}{45} + \frac{\pi^2}{12}\left(\frac{m^2}{T^2}\right) - \frac{\pi}{6}\left(\frac{m^2}{T^2}\right)^{3/2} - \frac{m^4}{32T^4}\ln\left(\frac{m^2}{a_BT^2}\right) + \cdots, \tag{16}$$

where $\log a_B = 5.4076$. Since the scalar ϕ obtains the thermal mass $m_{th}^2 \simeq T^2\lambda_7/6$, the finite temperature effective potential has a minimum at origin when the temperature of the Universe is high. It is easy to find that ϕ becomes tachyonic below the temperature

$$T_c \simeq \sqrt{\frac{6m_\phi^2}{\lambda_7}}, \tag{17}$$

and starts developing non-zero VEV. In this way, the Majorana masses of $\chi_{L,R}$ are generated after this phase transitions at a late-time.

In this paper we consider the case that $T_c < M_{\chi\psi}^{Dirac}/20$, which means the Dirac asymmetric dark matter particle becomes Majorana at a late-time well after the decoupling of the efficient annihilation processes of the dark matters and anti-dark matters. The DM mass matrix is now modified to include the Majorana mass terms such as

$$M = \begin{pmatrix} y\frac{v_s}{\sqrt{2}} & \mu_\chi & 0 & Y\frac{v}{\sqrt{2}} \\ \mu_\chi & y\frac{v_s}{\sqrt{2}} & Y\frac{v}{\sqrt{2}} & 0 \\ 0 & Y\frac{v}{\sqrt{2}} & 0 & \mu_\psi \\ Y\frac{v}{\sqrt{2}} & 0 & \mu_\psi & 0 \end{pmatrix}. \quad (18)$$

Therefore, the Dirac fermion asymmetric dark matter before the phase transition by $\langle\phi\rangle$ has taken the form of the lightest Majorana fermion mass eigenstate which is a linear combination among $\chi_{L,R}$ and $\psi_{L,R}$:

$$\chi_j = N_{j1}\chi_L^c + N_{j2}\chi_R + N_{j3}\psi_L^c + N_{j4}\psi_R, \quad (19)$$

where N_{ij} is the unitary matrix to diagonalize the mass matrix. When we take $\mu_\psi = \mu_\chi(=\mu)$ for example, the mass eigenvalues are given by

$$\frac{\sqrt{8Y^2v^2 + 2y^2v_s^2} + \sqrt{2}yv_s}{4} \pm 4\mu, \quad (20)$$

and the lighter one is the dark matter mass.

The doublet scalar η is heavier than the dark matter fermion and decays into the dark matter and the SM particles. The main decay mode is found to be $\eta \to \chi_j\bar{\nu}$ through the mixing between the heavy right-handed neutrinos and the SM neutrinos, $\sqrt{\frac{m_\nu}{M_N}}$, in the seesaw mechanism [42], where m_ν is the typical light neutrino mass scale. We estimate the decay width as

$$\Gamma(\eta \to \bar{\nu}\chi_j) \simeq \frac{1}{16\pi}Y_N^2 M_\eta \left(\frac{m_\nu}{M_N}\right), \quad (21)$$

which corresponds to

$$\tau(\eta \to \bar{\nu}\chi_j) \simeq 3.3 \times 10^{-5}\left(\frac{10^{-4}}{Y_N}\right)^2\left(\frac{0.1\text{eV}}{m_\nu}\right)\left(\frac{M_N}{M_\eta}\right)[\text{sec}]. \quad (22)$$

We can choose the parameters so as for the η to decay after the decoupling of pair annihilation processes of dark matters and also ηs. As we have discussed before, the N_R decay generates the same amounts for the dark matter asymmetry and the η asymmetry in the Universe. After the very efficient annihilation, only the dark matters are left in the Universe, in the same way, only the same amount of ηs exist. The late-time η decay creates the anti-dark matter and hence, the total dark matter number in the Universe becomes zero after the η decay. However, at that time, the dark matter pair annihilation processes have already frozen out, and the total number density of the dark matter (plus anti-dark matter) remains twice the number density of the dark matter just before the decay of the η.

Associated with the spontaneous global $U(1)$ symmetry breaking by $\langle\phi\rangle \neq 0$, cosmic strings would be formed. However, this is cosmologically harmless since the VEV scale of ϕ is small and the mass per unit length of cosmic strings is very small. In addition, associated with this phase transition the NG mode appears. However, this NG mode does not couple with SM particles, and the constraints from e.g., the cooling rate of stars through the NG mode emissions inside the core can be avoided. Since the NG mode was in thermal equilibrium, its energy density contributes to the energy density of the relativistic particles. We estimate the contribution to the extra neutrino species at the Big Bang Nucleosynthesis era $T_{BBN} \simeq 1$ MeV as

$$\Delta N_{\rm eff} \simeq \left.\frac{\rho_{\rm NG}}{\rho_\nu}\right|_{T_{BBN}} \approx 0.2. \tag{23}$$

This extra contribution is interesting in terms of the value $N_{\rm eff} = 3.85 \pm 0.62$ recently reported in Ref. [43] using the data from the South Pole Telescope.

C. Direct/Indirect detection of dark matter

There is a variety of ongoing and planned experiments to detect a dark matter particle directly or indirectly, through the elastic scattering of dark matter particle off with nuclei. The dark matter in our model couples with quarks in two ways. One is through Higgs bosons relevant to the spin-independent elastic scattering process, the other is the Z boson exchange which causes the spin-dependent elastic scattering of the dark matter particle. Note that because of its Majorana nature, the dark matter particle has only the axial vector coupling with the Z boson, and the Z boson exchange process has no contribution to the spin-independent elastic scattering process. If the dark matter is the Dirac fermion in the absence of ϕ VEV, the Z boson exchange dominantly contributes to the spin-independent elastic scattering process and as a result, our scenario will be excluded as the dark matter scenario with the heavy Dirac neutrino DM [25] or the left-handed sneutrino DM [44] [2]. Therefore, the late-time Majorana mass generation is crucial for our dark matter scenario to be phenomenologically viable.

The cross section of the spin-independent elastic scattering with a proton is given by

$$\sigma_{\rm SI}^{(p)} = \frac{4}{\pi}\left(\frac{m_p m_\chi}{m_p + m_\chi}\right)^2 f_p^2, \tag{24}$$

with the hadronic matrix element

$$\frac{f_p}{m_p} = \sum_{q=u,d,s} f_{Tq}^{(p)}\frac{\alpha_q}{m_q} + \frac{2}{27} f_{TG}^{(p)} \sum_{c,b,t}\frac{\alpha_q}{m_q}, \tag{25}$$

where m_q is a mass of a quark, and $f_{Tq}^{(p)}$ and $f_{TG}^{(p)}$ are constants. α_q is the coefficient of the effective operator $\mathcal{L}_{int} = \alpha_q \chi_j \chi_j \bar{q} q$ for this scattering process and is given by

$$\alpha_q = \frac{1}{2}\sum_i \frac{(N_{j1}^2 + N_{j2}^2) y S_{i1} + 2(N_{j2}N_{j3} + N_{j1}N_{j4}) Y S_{i2}}{\sqrt{2} v M_{H_i}^2} m_q S_{i2}. \tag{26}$$

For $y v_s \ll Y v$, α_q is simplified as

$$\begin{aligned} \alpha_q &\simeq \sum_i \frac{y S_{i1} \pm Y S_{i2}}{4\sqrt{2} v M_{H_i}^2} m_q S_{i2} \\ &= \frac{m_q}{4\sqrt{2} v}\left(\frac{y\cos\alpha \pm Y \sin\alpha}{M_{H_1}^2}\sin\alpha + \frac{-y\sin\alpha \pm Y\cos\alpha}{M_{H_2}^2}\cos\alpha\right), \end{aligned} \tag{27}$$

where $\pm$ corresponds to which mass eigenvalue in Eq. (20) is the dark matter mass.

When the Yukawa couplings are non-universal,

$$\mathcal{L}_{Yukawa} \supset Y_1 \bar{\psi}_L \Phi \chi_R + Y_2 \bar{\chi}_L \Phi \psi_R + h.c., \tag{28}$$

we also consider the spin-dependent scattering processes of the dark matter. The corresponding cross section with N-nucleus of mass m_N and the spin J is given by

$$\sigma_{\rm SD} = \frac{32}{\pi} G_F^2 \left(\frac{m m_N}{m + m_N}\right)^2 \Lambda^2 J(J+1), \tag{29}$$

with

$$\Lambda = \frac{1}{J}(a_{\rm p}\langle S_{\rm p}\rangle + a_{\rm n}\langle S_{\rm n}\rangle), \tag{30}$$

[2] Mixed [12, 45] or some right-handed [46, 47] sneutrinos are still viable WIMP DM candidates.

and the effective coupling with a proton and a neutron

$$a_{\rm p} = \sum_{q=u,d,s} \frac{\alpha_2}{\sqrt{2}G_F} \Delta_q^{(p)}, \tag{31}$$

$$a_{\rm n} = \sum_{q=u,d,s} \frac{\alpha_2}{\sqrt{2}G_F} \Delta_q^{(n)}, \tag{32}$$

where G_F is the Fermi constant, $\langle S_{\rm p(n)} \rangle$ is the average of spin of proton (neutron) in a nuclei, and Δ_q denotes the quark spin content. α_2 is the coefficient of the effective operator $\mathcal{L}_{int} = \alpha_2 \chi_j \gamma_5 \chi_j \bar{q} \gamma_5 q$ mediated by the Z boson, and this is proportional to $N_{j3}^2 - N_{j4}^2$. For $yv_s \ll Yv$, this is simplified as

$$N_{j3}^2 - N_{j4}^2 \simeq \frac{(Y_1^2 - Y_2^2)}{2\sqrt{(Y_1^2 - Y_2^2)^2 + 8(Y_1 + Y_2)^2 (\mu/v)^2}}. \tag{33}$$

D. Benchmarks

There are a lot of parameters in our model and we here choose typical parameter sets for our benchmark scenarios. In order to realize our scenario, very efficient annihilation processes of the asymmetric dark matters are necessary. In addition, it is interesting if there exists a parameter regions which leads to the spin-independent/dependent elastic scattering cross sections testable by the future experiments.

1. $M_{\chi\psi}^{Dirac} > M_W$ *and* $\alpha \simeq 0$

Firstly, we consider the case where $M_{\chi\psi}^{Dirac} > M_W$ and the mixing angle of Higgs bosons is negligible, $\alpha \simeq 0$ [3]. In this case, the Dirac fermion dark matters (and also the η scalars) sufficiently annihilate into mainly W^+W^- through the $SU(2)$ gauge interaction in the early Universe. After the phase transition by non-zero $\langle \phi \rangle$, the dark matter becomes Majorana but its relic number density has been already determined by the dark matter asymmetry at the period when that particle was Dirac. At present, since this dark matter particle is Majorana, the cross section with nuclei which is relevant to the dark matter direct detection is induced by not the Z boson exchange but Higgs boson exchanges. For $\alpha \simeq 0$, the SM-like

[3] Since we have fixed $v_s \ll v$, $\alpha \ll 1$ is natural for $\lambda_1 \sim \lambda_4 \sim \lambda_7$ from Eq. (5).

Higgs (H_2) exchange process is dominant and we can use the simplified form of α_q as

$$\alpha_q \simeq \frac{m_q}{4\sqrt{2}v}\frac{Y}{M_{H_2}^2}. \tag{34}$$

For instance, for $M_{H_2} \simeq 120$ GeV and $Y \simeq 0.2$, we find

$$\sigma_{\rm SI}^{(p)} \simeq 10^{-9}{\rm pb}, \tag{35}$$

which is close to the current XENON 100 exclusion limit of $\sigma_{\rm SI}^{(p)} \gtrsim 10^{-8}$ pb [48].

Similarly, we estimate the spin-dependent cross section with a proton. For instance, for $\mu \simeq 100$ GeV and $(Y_1, Y_2) \simeq (0.4, 0.2)$, we obtain

$$\sigma_{\rm SD}^{(p)} \simeq 3 \times 10^{-4}{\rm pb}, \tag{36}$$

almost independently of the dark matter mass. The terrestrial experiments of dark matter direct search give the upper bound on $\sigma_{\rm SD}^{(p)} <$ several $\times 10^{-2}$ pb by the SIMPLE [49]. In fact, the most stringent upper bound is derived from the indirect dark matter search by IceCube [50] as $\sigma_{\rm SD}^{(p)} \lesssim 10^{-3} - 10^{-4}$pb, depending on dark matter mass. Our result is close to this IceCube bound.

2. $M_{\chi\psi}^{Dirac} < M_W$ and $\alpha \simeq 0$

Next, we consider the case where $M_{\chi\psi}^{Dirac} < M_W$ and the mixing angle $\alpha \simeq 0$. Since the annihilation channel to W^+W^- is kinematically forbidden in this case, the Dirac fermion dark matters mainly annihilate to a light ϕ pair through the Yukawa coupling. Since we have assumed the mass spectrum $M_\eta > M_{\chi\psi}^{Dirac}$, even in this case, we may expect the annihilation of η into W-boson pairs. The annihilation cross section of Dirac fermion dark matters is estimated as

$$\sigma v \simeq \frac{9(NyN)^4}{16\pi m_\chi^2}\frac{T}{m_\chi}, \tag{37}$$

which can be large enough for $y = 0.1 - 1$. The spin-independent/dependent cross sections can be the same as those in the case with $M_{\chi\psi}^{Dirac} > M_W$.

3. Sizable α

Finally we consider the case with a sizable α. In this case, the cross section with a proton can be enhanced by the first term in Eq. (27) because the singlet-like Higgs (H_1) is light. A portion of parameter space with very light H_1 and/or not small y can be excluded by the current bound from the dark matter direct search experiments.

E. Precision measurement

In our scenario, the dark matter particle is a mixture of the $SU(2)$ doublet and singlet fermions after the electroweak symmetry breaking, which causes the iso-spin violation, in other words, the mass splitting between up-sector and down-sector fermions in the $SU(2)$ doublet. Such a mass splitting induces additional contribution to the ρ parameter, whose deviation from 1 is very severely constrained, $\Delta\rho = \rho - 1 = \mathcal{O}(10^{-3})$ [51]. Assuming a small Majorana mass term, $yv_s/\sqrt{2}$, we evaluate 1-loop corrections to $\Delta\rho$ via the dark matter fermions with the mass matrix in Eq (11). For $\mu_\chi > Yv/\sqrt{2}$, we find

$$\Delta\rho \simeq \frac{Y^2}{4\pi^2}\left(\frac{\mu_\psi}{\mu_\chi - \mu_\psi}\right). \tag{38}$$

The constraint is satisfied for $Y \lesssim 0.3$ with $\mu_\chi \gtrsim 3\mu_\psi$. Thus, our discussion on the direct/indirect detection of dark matter is consistent with the ρ parameter constraint.

III. MODEL B

A. Setup

Next we consider another realization of our scenario in the context of a gauged $U(1)_{B-L}$ extended model. In addition to the minimal gauged $U(1)_{B-L}$ extension with the right-handed neutrinos N_R and $B-L$ Higgs field Ψ [52–57], we further introduce a vector-like pair of SM singlet fermions ($\psi_{L,R}$) with $B-L$ charge -2, another vector-like pair of totally singlet fermions ($\chi_{L,R}$) and two $B-L$ charged/uncharged complex scalars (η and ϕ). The charge assignments are shown in Table II. Here we have introduced a discrete Z_2 parity to ensure the stability of dark matter particle and a global $U(1)$ symmetry to forbid Majorana mass terms for the $B-L$ charge neutral fermions.

Lagrangian relevant for our discussion is given by

$$\begin{aligned} -\mathcal{L} \supset\ & \frac{1}{2}y\bar{\chi_R^c}\phi\chi_R + \frac{1}{2}y\bar{\chi_L^c}\phi\chi_L + \frac{1}{2}\bar{N_R^c}(Y_\Psi\Psi^\dagger)N_R \\ & +\mu_\chi\bar{\chi_L}\chi_R + \mu_\psi\bar{\psi_L}\psi_R + Y\bar{\psi_L}\Psi\chi_R + Y\bar{\chi_L}\Psi^\dagger\psi_R + Y_N\bar{\psi_L}\eta^\dagger N_R + h.c. \\ & +V(\Phi,\Psi,\eta,\phi) \end{aligned} \tag{39}$$

TABLE II: Particle contents of Model B

Fields	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$	Z_2-parity	Global U(1)
ψ_L	1	0	-2	$-$	-1
ψ_R	1	0	-2	$-$	-1
χ_L	1	0	0	$-$	-1
χ_R	1	0	0	$-$	-1
Ψ	1	0	-2	$+$	0
η	1	0	$+1$	$-$	$+1$
ϕ	1	0	0	$+$	$+2$
N_R	1	0	-1	$+$	0

with

$$
\begin{aligned}
V(\Phi,\Psi,\eta,\phi) &= m_1^2|\Phi|^2 + m_2^2|\Psi|^2 + \lambda_1|\Phi|^4 + \lambda_2|\Psi|^4 + \lambda_3|\Phi|^2|\Psi|^2 \\
&\quad + m_3^2|\phi|^2 + (\lambda_4|\Phi|^2 + \lambda_5|\Psi|^2 + \lambda_6|\eta|^2)|\phi|^2 + \lambda_7|\phi|^4 \\
&\quad + m_4^2|\eta|^2 + (\lambda_8|\Phi|^2 + \lambda_9|\Psi|^2)|\eta|^2 + \lambda_{10}|\eta|^4 \\
&\quad + \lambda_W \eta\eta\Psi\phi^\dagger + h.c.,
\end{aligned}
\tag{40}
$$

where Φ denotes the SM Higgs doublet field. For simplicity, we have assumed the universal Yukawa coupling constants y and Y. We fix the parameters in the scalar potential so as for Φ, Ψ and ϕ to develop VEVs and in this case the physical Higgs bosons are given by

$$
\Phi = \begin{pmatrix} 0 \\ \frac{1}{\sqrt{2}}(v+h) \end{pmatrix}, \ \Psi = \frac{1}{\sqrt{2}}(v'+H), \ \phi = \frac{1}{\sqrt{2}}(v_s+\varphi) \tag{41}
$$

around their VEVs (v, v' and v_s). The mass eigenstates are obtained by the unitary rotation

$$
\begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} = \begin{pmatrix} S_{11} & S_{12} & S_{13} \\ S_{21} & S_{22} & S_{23} \\ S_{31} & S_{32} & S_{33} \end{pmatrix} \begin{pmatrix} \varphi \\ h \\ H \end{pmatrix}, \tag{42}
$$

by which the mass matrix

$$\text{Mass term } = \frac{1}{2}(\varphi\ h\ H)\begin{pmatrix} 2v_s^2\lambda_7 & vv_s\lambda_4 & v'v_s\lambda_5 \\ vv_s\lambda_4 & 2v^2\lambda_1 & vv'\lambda_3 \\ v'v_s\lambda_5 & vv'\lambda_3 & 2v'^2\lambda_2 \end{pmatrix}\begin{pmatrix} \varphi \\ h \\ H \end{pmatrix} \tag{43}$$

is diagonalized.

B. Cosmological evolution

When Ψ develops the VEV $\langle\Psi\rangle = \frac{v'}{\sqrt{2}}$, the $B-L$ gauge symmetry is broken. Associated with the symmetry breaking, the $B-L$ gauge boson (Z') and the right-handed neutrinos acquire their masses, $m_{Z'}^2 = 4g_{B-L}^2 v'^2$ and $m_{N_R} = \frac{Y_\Psi v'}{\sqrt{2}}$. We set this symmetry breaking at a energy much higher than the electroweak symmetry breaking. The current lower bound on v' is found to be $v' \gtrsim 3$ TeV [58, 59], and our setup is consistent with it.

Right-handed neutrinos have two decay modes; decays into lepton and Higgs ($N_R \to L\Phi$) and into η and ψ ($N_R \to \psi\eta$). As the same as the discussion for Model A in the previous section, the CP-asymmetric decays of the right-handed neutrinos generate the lepton asymmetry and the asymmetry of ψ and η. We assume the scalar η is the heavier than the other Z_2-odd particles. The ψ-asymmetry generated by N_R decay is nothing but the asymmetry of the dark matter sector. Notice that the same amount of asymmetry between η and η^* is also generated as in Model A.

After the $B-L$ symmetry is broken by $\langle\Phi\rangle = v'/\sqrt{2}$ (but ϕ has not developed the VEV at this moment), Z_2-odd Dirac fermion mass matrix is given by

$$(\bar{\chi_L}\bar{\chi_R^c}\bar{\psi_L}\bar{\eta_R^c})M\begin{pmatrix} \chi_L^c \\ \chi_R \\ \psi_L^c \\ \psi_R \end{pmatrix} \tag{44}$$

with

$$M = \begin{pmatrix} y\langle\phi\rangle & \mu_\chi & 0 & Y\frac{v'}{\sqrt{2}} \\ \mu_\chi & y\langle\phi\rangle & Y\frac{v'}{\sqrt{2}} & 0 \\ 0 & Y\frac{v'}{\sqrt{2}} & 0 & \mu_\psi \\ Y\frac{v'}{\sqrt{2}} & 0 & \mu_\psi & 0 \end{pmatrix}, \tag{45}$$

where $\langle\phi\rangle = 0$. The mass eigenstates are mixture of two Dirac fermions, χ and ψ and their mass eigenvalues are given by

$$M_{\chi\psi}^{Dirac} = \frac{\mu_\chi + \mu_\psi \pm \sqrt{2Y^2 v'^2 + (\mu_\chi - \mu_\psi)^2}}{2}. \tag{46}$$

The lighter mass eigenstate is nothing but the asymmetric Dirac fermion dark matter.

The WIMP dark matter scenario in the context of the gauged $B - L$ model has been investigated before, and it has been shown that the dark matter annihilation processes are not efficient. Only in the limited cases with the Higgs boson resonance [60–62] or the the Z' boson resonance [62, 63], the relic abundance of the dark matter particle can be lower than the observed value. Therefore, in order to realize our asymmetric dark matter scenario, we need to tune the dark matter masses $M_{\chi\psi}^{Dirac}$ to be close to half of either Higgs bosons or Z' boson masses. So is the η mass, in order to realize efficient annihilations between η and η^*.

At a low energy where only ϕ is light and the others are decoupled, the scalar potential is simplified as

$$V(\Phi, \eta, \phi) \simeq -m_\phi^2 |\phi|^2 + \lambda_7 |\phi|^4, \tag{47}$$

where $m_\phi^2 = -m_3^2 - \frac{\lambda_4 v^2 + \lambda_5 v'^2}{2} > 0$, and the scalar ϕ develops the VEV. In the early Universe, we consider the thermal effects on the scalar potential. By the same analysis in the previous section, the thermal mass term for ϕ is found to be $m_{th}^2 \simeq T^2 \lambda_7/6$ in the high-temperature expansions. Thus, the phase transition occurs at the temperature $T_c \simeq \sqrt{6m_\phi^2/\lambda_7}$. Again, we assume a small m_ϕ and this phase transition takes place well after the dark matter and anti-dark matter annihilation have frozen out. The same arguments on cosmic strings and NG mode for Model A are applicable to this model.

As the same as in the Model A, the η decay produce anti-dark matter at late-time. Since this happens after the freeze out of the annihilation processes of dark matters and ηs, the produced anti-dark matters remain without annihilations with the same number density as the dark matter one.

C. Direct dark matter detection

Once the ϕ has developed the VEV in a late-time, the fermion mass matrix in Eq. (45) has non-zero diagonal elements and the dark matter particle has taken the form of the

Majorana fermion. Because of the Majorana nature, the dark matter particle has the axial vector coupling with the Z' boson, while the SM quarks have the vector coupling. As a result, there is no effective couplings between the dark matter and the SM quarks induced by the Z' boson exchange in the non-relativistic limit. Therefore, for Model B, we only consider the spin-independent elastic scattering process via Higgs boson exchanges.

The coefficient of the operator $\mathcal{L}_{int} = \alpha_q \chi_j \chi_j \bar{q} q$ is found to be

$$\alpha_q = \frac{1}{2} \sum_i \frac{(N_{j1}^2 + N_{j2}^2) y S_{i1} + 2(N_{j2}N_{j3} + N_{j1}N_{j4}) Y S_{i3}}{\sqrt{2} v M_{H_i}^2} m_q S_{i2}, \tag{48}$$

where N_{ij} is the unitary matrix to diagonalize the mass matrix. In the limit $y v_s \ll Y v'$ for simplicity, α_q is reduced as

$$\alpha_q \simeq \frac{1}{2} \sum_i \frac{y S_{i1} \pm 2 Y S_{i3}}{2\sqrt{2} v M_{H_i}^2} m_q S_{i2}. \tag{49}$$

D. Benchmark scenarios

As we have discussed above, viable parameter sets are very limited in order to achieve the large annihilation cross section for dark matter particles via the enhancement of the Higgs or Z' boson resonances. Let us consider two benchmarks.

1. The Z' boson resonance

We take the asymmetric dark matter mass to be $M_{\chi\psi}^{Dirac} \simeq M_{Z'}/2$, so that the dark matter annihilation cross section is enhanced [62, 63] and the relic dark matter originates from the dark matter asymmetry in the Universe. Similarly, we take the η mass to be around half of the heavy Higgs boson mass. In this case there is little correlation between the annihilation cross section and the spin-independent cross section of the dark matter off the nuclei, except the dark matter mass being half of Z' boson mass.

2. A Higgs boson resonance

Another benchmark is to fix the dark matter mass $M_{\chi\psi}^{Dirac} \simeq M_{H_i}/2$ and $M_\eta \simeq M_{H_j}/2$ or $M_{Z'}/2$, where the Dirac fermion annihilation cross section is enhanced by the s-channel Higgs resonances and η annihilation cross section is enhanced by the s-channel Higgs or Z'

resonances as discussed in Refs. [60, 61] for similar models. For simplicity, we fix the dark matter mass to be half of the SM-like Higgs boson and in this case the structure of our model is similar to the so-called Higgs portal dark matter scenario. A representative model is the SM plus a gauge singlet real scalar as the dark matter [64–67]. Since the sensitivity of the dark matter direct detection experiments has been greatly improved, the allowed parameter region of the model to simultaneously satisfy the constraints from the relic abundance and the direct detection is very limited [68–71]: For relatively light dark matter, say, $M_{DM} \lesssim 1$ TeV, the dark matter mass should be around half of the Higgs boson mass ($M_{DM} \approx M_H/2$). The correct relic abundance is achieved by the Higgs resonance even though the coupling between the dark matter and the Higgs boson is very small. On the other hand, such a small coupling makes the direct detection of the dark matter very hard. If the dark matter is lighter than half of the Higgs mass and the dark matter mass is not just on the Higgs pole, we can find the Higgs portal dark matter signal at high energy colliders such as the LHC [70, 71] through the invisible decay of the Higgs boson into the dark matter particle pair [72, 73].

In our model, there is a crucial difference from the Higgs portal dark matter. For the asymmetric dark matter scenario, the large annihilation cross section is welcome in order to completely erase the symmetric part of the relic abundance of dark matter and anti-dark matter, leaving only the asymmetric part. Therefore, it is not necessary for the coupling between the dark matter and the Higgs boson to be small. In this case the spin-independent cross section can be accessible to the future experiments.

IV. SUMMARY AND DISCUSSIONS

We have proposed a novel asymmetric dark matter scenario. The dark matter particle is initially the Dirac fermion and the dark matter asymmetry is generated by the same mechanism as leptogenesis. Through efficient dark matter annihilation processes, the symmetric part of the dark matter number density is erased and the dark matter relic density is determined by the dark matter asymmetry. After the annihilation processes have frozen out, a phase transition of a light scalar field takes place and generates Majorana mass terms for the dark matter particles. As a result, the originally asymmetric dark matter turns into the Majorana fermion. Although the dark matter behaves as the WIMP at present, its relic abundance is basically independent of the pair annihilation cross section.

In order to address this scenario in detail, we have proposed two simple models. The first model is based on the SM gauge group and no new gauge interaction is introduced. The dark matter originates from the $SU(2)_L$ doublet Dirac fermion and similar to the Higgsino in the MSSM. Through the coupling with the right-handed neutrinos, the dark matter asymmetry in the Universe is generated by the out-of-equilibrium decay of the right-handed neutrinos, the same mechanism as leptogenesis. The annihilation cross section through the $SU(2)_L$ gauge interaction is large enough to erase the symmetric part of the dark matter and anti-dark matter abundance. At a low energy after the freeze-out of the annihilation process, the thermal effects for the scalar potential becomes smaller, and a light scalar field develops the VEV. The Majorana mass term generated by the VEV turns the dark matter particle into the Majorana particle. Because of this Majorana nature, the dark matter has no spin-independent interaction with the quarks mediated by the Z boson. We have estimated the spin independent and independent elastic scattering cross sections of the dark matter with a proton and found that the cross sections can be close to the current experimental bounds with reasonable model-parameter choices. Therefore, the dark matter can be detected in the near future.

In the second model, we have introduced an extra $U(1)_{B-L}$ gauge interaction and all new particles introduced are singlet under the SM gauge group. Efficient Dirac fermion dark matter annihilation is achieved through the resonance effect by the Higgs or Z' boson mediated processes, as a result, the dark matter mass should be around half of either Higgs or Z' boson masses. As in the first model, the phase transition of a light scalar occurs after the annihilation processes have been frozen out, and the dark matter takes the form of the Majorana fermion at a late-time. This second model is similar to the Higgs portal dark matter scenario. A crucial difference is that there is no constraint on the coupling between the dark matter and Higgs bosons from the dark matter relic abundance, because the relic abundance is determined by the dark matter asymmetry. Therefore, the spin-independent cross section can be as large as the current upper limit from the direct dark matter search experiments, independently of the annihilation cross section. This situation is similar to the non-thermal dark matter scenario, where the dark matter particle is produced by unstable relics such as moduli [74], Q-ball [75] and axino [76].

Finally, we note that the realization of our scenario in the context of MSSM is challenging because the Majorana mass terms for the gauginos is always present. Recently, it has been shown [77] that an extreme parameter choice can realize the Higgsino to be a viable asymmetric dark matter candidate.

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Appendix A

Here, we note the explicit form of N.

$$N = \begin{pmatrix} -\frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_+} & \frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_+} & -\frac{1}{2}D_+ & \frac{1}{2}D_+ \\ \frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_-} & -\frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_-} & -\frac{1}{2}D_- & \frac{1}{2}D_- \\ -\frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_+} & -\frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_+} & \frac{1}{2}D_+ & \frac{1}{2}D_+ \\ \frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_-} & \frac{Yv}{\sqrt{(2Yv)^2+(yv_s)^2}D_-} & \frac{1}{2}D_- & \frac{1}{2}D_- \end{pmatrix}, \tag{A1}$$

with

$$D_+ = \sqrt{1+\frac{yv_s}{\sqrt{(2Yv)^2+(yv_s)^2}}}, \tag{A2}$$

$$D_- = \sqrt{1-\frac{yv_s}{\sqrt{(2Yv)^2+(yv_s)^2}}}. \tag{A3}$$

Appendix B: Amplitude

We give explicit formulas of the invariant amplitude squared for the pair annihilation of Dirac dark matter χ into light singlet scalar ϕ pair.

$$w(s;\to\phi\phi) \equiv \frac{1}{4}\int d\mathrm{LIPS}\overline{|\mathcal{M}|}^2 \simeq \frac{1}{32\pi}\left(NyN\right)^4\left(F(s)+G(s)\ln\left|\frac{a+b}{a-b}\right|\right), \tag{B1}$$

The auxiliary functions appear above are defined as

$$F(s) \equiv 2\left(-2 - \frac{(4m_\chi^2 - m_\phi^2)^2}{m_\chi^2(s - 4m_\phi^2) + m_\phi^4}\right), \tag{B2}$$

$$G(s) \equiv 2\sqrt{\frac{s - 4m_\phi^2}{s - 4m_\chi^2}}\frac{\left(-32m_\chi^4 - 4m_\phi^2(s + 4m_\chi^2) + 16sm_\chi^2 + 6m_\phi^4 + s^2\right)}{s^2 - 6sm_\phi^2 + 8m_\phi^4}, \tag{B3}$$

$$a(s) \equiv s - 2m_\phi^2, \tag{B4}$$

$$b(s) \equiv \sqrt{(s - 4m_\phi^2)(s - 4m_\chi^2)}. \tag{B5}$$

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