



Baryonic and leptonic GeV dark matter

Bartosz Fornal^{*}, Alec Hewitt, Yue Zhao

Department of Physics and Astronomy, University of Utah, Salt Lake City, UT 84112, USA

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ABSTRACT

We perform a systematic analysis of models with GeV-scale dark matter coupled to baryons and leptons. Such theories provide a natural framework to explain the matter-antimatter asymmetry of the universe. We find that only a few baryonic dark matter models are free from tree-level proton decay without explicitly imposing baryon number conservation, provided that the dark matter particle is sufficiently heavy. We enumerate those cases and present a brief overview of their phenomenology. We then focus on a leptonic dark matter model for a more detailed discussion of the baryon asymmetry generation via leptogenesis, the symmetry restoration in the dark sector and the expected dark matter annihilation signals in indirect detection experiments.

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1. Introduction

The evidence for dark matter in the universe is indisputable. Not only its existence, but also its distribution and abundance are precisely determined from various measurements, including galaxy rotation curves [1], cosmic microwave background [2], gravitational lensing [3], etc. Nevertheless, the dark matter mass and its interactions with Standard Model particles remain a mystery. It is not even known if dark matter consists of elementary particles or macroscopic objects. The mass of an elementary dark matter particle can be anywhere between $\sim 10^{-31}$ GeV (fuzzy dark matter) [4,5] and $\sim 10^{19}$ GeV (WIMPzillas) [6,7] (despite the unitarity bound [8]), while the mass of macroscopic dark matter objects ranges from $\sim 10^{17}$ GeV (dark quark nuggets) [9] to $\sim 10^{59}$ GeV (primordial black holes) [10,11]. In most cases, the dark matter interactions with the known particles are small; from a theoretical perspective such interactions can even be absent.

Interestingly, the ratio of the abundances of dark matter and ordinary matter is on the order of five. This suggests that the two sectors may be related and, perhaps, share a common origin. This is precisely the idea behind theories of asymmetric dark matter [12,13], in which an effective interaction between dark matter and Standard Model particles is established. The explanation of the matter-antimatter asymmetry of the universe in those theories relies on the asymmetries in the two sectors being generated simultaneously and the dark matter particles being at the GeV scale.

In this letter, we systematically analyze scenarios in which dark matter couples to quarks and/or leptons. We study the possible effective operators describing such interactions up to dimension eight and analyze their particle model realizations. For a successful baryogenesis or leptogenesis, those operators must have a nonzero baryon number contribution carried by the quarks or a nonzero lepton number from the leptons. The effective dark matter-Standard Model interactions require either scalar or vector mediators to be present at the particle level of each model. For dark matter coupled to baryons (baryonic dark matter) at least one mediator in each model is necessarily a color triplet. For dark matter coupled to leptons (leptonic dark matter) the mediators do not carry color.

We demonstrate that for baryonic and baryoleptonic dark matter the possible mediators are: the scalars $(3, 1)_{-4/3}$, $(3, 1)_{-1/3}$, $(3, 1)_{2/3}$, or the vectors $(3, 2)_{-5/6}$, $(3, 2)_{1/6}$. Without imposing an additional symmetry, all of those particles, except for the scalar $(3, 1)_{2/3}$, can trigger tree-level proton decay to Standard Model particles [14,15]. Thus the mass of those mediators is elevated above $\sim 10^{16}$ GeV by the stringent experimental constraints on the proton lifetime, limiting their capabilities of explaining baryogenesis and their experimental probes.

Such large mediator masses can be avoided by fine-tuning some of the couplings to be small or by explicitly imposing baryon and/or lepton number conservation. However, there is no strong theoretical argument to expect baryon or lepton number to be exact symmetries of nature. In fact, both of them are already violated within the Standard Model itself at the non-perturbative level by the electroweak sphalerons. Guided by the requirement of proton stability without assuming a fine-tuning of couplings or imposing an additional symmetry, we focus on baryonic dark matter models

^{*} Corresponding author.

E-mail addresses: bfornal@physics.utah.edu (B. Fornal), u0961360@uemail.utah.edu (A. Hewitt), zhaoyue@physics.utah.edu (Y. Zhao).

involving only the scalar mediator $(3, 1)_{2/3}$, as well as on leptonic dark matter models. The latter are naturally free from tree-level proton decay, since they involve the color singlet scalar mediators $(1, 1)_1$ or $(1, 2)_{-1/2}$.

We find that only some of the baryonic dark matter direct detection signatures considered in the literature are realized in models that do not suffer from tree-level proton decay to Standard Model particles. In particular, the dark matter-nucleon annihilation, which is possible in such models, necessarily involves a kaon in the final state. This conclusion increases the importance of the Deep Underground Neutrino Experiment (DUNE) [16] in complementing the efforts of Super-Kamiokande [17] in the search for baryonic dark matter. Regarding leptonic dark matter models, based on a concrete example, we analyze their potential for generating the matter-antimatter asymmetry through leptogenesis and highlight their unique feature of symmetry restoration in the dark sector, achieved when one of the dark sector particles is unstable on long time scales and decays to dark matter. Such a rebirth of symmetric dark matter leads to the possibility of enhanced annihilation signals at present times, previously considered in the context of heavy dark matter [18–20] and oscillating dark matter [21–25], and results in signatures that can be searched for by the Fermi Gamma-Ray Space Telescope [26] and the future e-ASTROGAM [27].

2. The models

In this section, we enumerate the possible particle physics models for the effective operators describing the interactions of dark matter with just the quarks (baryonic dark matter), both quarks and leptons (baryoleptonic dark matter), and with leptons only (leptonic dark matter), and briefly discuss their phenomenology. We adopt the four-component Dirac spinor notation for the fermion fields, indicating with a subscript L or R their left- or right-handed chirality. The conjugate fields are denoted by a bar symbol. We focus on the operators which carry a nonzero baryon number contribution from quarks or a nonzero lepton number contribution from leptons, since only those theories provide promising asymmetric dark matter frameworks. In particular, we do not consider operators of the type $q\bar{q}\chi\bar{\chi}$ or $\ell\bar{\ell}\chi\bar{\chi}$, which arise, e.g., in certain theories with gauge bosons coupled to baryon or lepton number [28,29]. In the following analysis, χ and $\tilde{\chi}$ are Dirac fermions, whereas ϕ and Φ are complex scalars. Both χ and ϕ are Standard Model singlets.

Baryonic dark matter

The simplest effective operator describing the interaction of dark matter with quarks is the dimension six $qqq\chi$, where q stands for Q_L , d_R or u_R . The possible gauge-invariant realizations are $u_R d_R d_R \chi$ and $Q_L Q_L d_R \chi$. There are three models one can write down for those operators, labeled as Models 1A–1C in Table 1. They involve the mediators: scalar $(3, 1)_{2/3}$, scalar $(3, 1)_{-1/3}$, and vector $(3, 2)_{1/6}$, respectively. Among those particles, only the scalar $(3, 1)_{2/3}$ in Model 1A does not give rise to tree-level proton decay to Standard Model particles, since the quantum numbers do not allow Φ to couple to a quark and a lepton [14,15]. This is in contrast to the other two cases, where the scalar $(3, 1)_{-1/3}$ can have the couplings $\Phi u_R d_R$ and $\Phi^* u_R e_R$, whereas the vector $(3, 2)_{1/6}$ can couple via $X_\mu Q_L \gamma^\mu d_R$ and $X_\mu^\dagger L_L \gamma^\mu u_R$, both leading to $p \rightarrow e^+ \pi^0$. For this reason, we focus on Model 1A below. An example of a diagram generating the operator $qqq\chi$ is shown in Fig. 1.

At dimension seven, the possible effective interactions are $qqq\chi\phi$ and $qqqH\chi$, where H is the Higgs field. In this study, we do not consider operators involving the Higgs field and only focus on the ones in the first category. There are six particle models for

Table 1

Effective operators describing the interaction of baryonic dark matter with quarks and their model realizations.

Model	Baryonic dark matter	
	Interactions	Mediators
$qqq\chi$		
1A	$\Phi d_R d_R, \Phi^* u_R \chi$	$\Phi = (3, 1)_{\frac{2}{3}}$
1B	$(\Phi u_R d_R \text{ or } \Phi Q_L Q_L), \Phi^* d_R \chi$	$\Phi = (3, 1)_{-\frac{1}{3}}$
1C	$X_\mu Q_L \gamma^\mu d_R, X_\mu^\dagger Q_L \gamma^\mu \chi$	$X_\mu = (3, 2)_{\frac{1}{6}}$
$qqq\chi\phi$		
2A	$\Phi d_R d_R, \Phi^* u_R \tilde{\chi}, \tilde{\chi} \chi \phi$	$\Phi = (3, 1)_{\frac{2}{3}}, \tilde{\chi} = (1, 1)_0$
2B	$\Phi d_R d_R, \Phi^* \tilde{\chi} \chi, \tilde{\chi} u_R \phi$	$\Phi = (3, 1)_{\frac{2}{3}}, \tilde{\chi} = (3, 1)_{\frac{2}{3}}$
2C	$(\Phi u_R d_R \text{ or } \Phi Q_L Q_L), \Phi^* d_R \tilde{\chi}, \tilde{\chi} \chi \phi$	$\Phi = (3, 1)_{-\frac{1}{3}}, \tilde{\chi} = (1, 1)_0$
2D	$(\Phi u_R d_R \text{ or } \Phi Q_L Q_L), \Phi^* \tilde{\chi} \chi, \tilde{\chi} d_R \phi$	$\Phi = (3, 1)_{-\frac{1}{3}}, \tilde{\chi} = (3, 1)_{-\frac{1}{3}}$
2E	$X_\mu Q_L \gamma^\mu d_R, \tilde{\chi} \chi \phi, X_\mu^\dagger Q_L \gamma^\mu \tilde{\chi}$	$X_\mu = (3, 2)_{\frac{1}{6}}, \tilde{\chi} = (1, 1)_0$
2F	$X_\mu Q_L \gamma^\mu d_R, X^\dagger \tilde{\chi} \chi, \tilde{\chi} Q_L \phi$	$X_\mu = (3, 2)_{\frac{1}{6}}, \tilde{\chi} = (3, 2)_{\frac{1}{6}}$

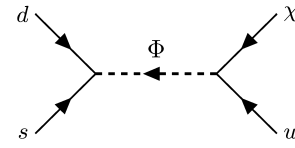


Fig. 1. Model 1A realization of the operator $qqq\chi$.

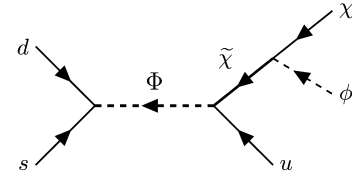


Fig. 2. Model 2A realization of the operator $qqq\chi\phi$.

the operator $qqq\chi\phi$, denoted as Models 2A–2F in Table 1. They involve the same scalar and vector mediators as introduced for the operator $qqq\chi$. In addition, an intermediate fermion $\tilde{\chi}$ is required (see, Fig. 2). We discuss in more detail Model 2A, since it does not suffer from tree-level proton decay to Standard Model particles.

At dimension eight, the operators are $qqq\chi\phi^2$, $qqqH\chi\phi$ and $qqqH^2\chi$. Again, we only consider the ones in the first category. Models for the operator $qqq\chi\phi^2$ can be constructed by introducing an additional scalar singlet field ϕ_e , replacing ϕ with ϕ_e in Models 2A–2F, and adding the interaction $\phi_e^* \phi^2$. The particle ϕ is then automatically stable without imposing specific relations between the masses. The particle χ , on the other hand, may be unstable if sufficiently heavy.

Model 1A ($qqq\chi$)

The Lagrangian for Model 1A is given by

$$-\mathcal{L}_1 \supset \lambda_q^{ab} \epsilon^{ijk} \Phi_i d_{Rj}^a d_{Rk}^b + \lambda_\chi^a \Phi^* \chi u_{Ri}^a + \text{h.c.}, \quad (1)$$

where i, j, k are color indices and a, b are flavor indices. Due to the antisymmetric nature of the ϵ tensor, the coupling λ_q^{ab} must be antisymmetric in flavor.

Although there is no tree-level proton decay to a final state consisting of only Standard Model particles, the model still suffers from proton decay if $m_\chi < m_p - m_e = 937.761$ MeV, since the proton can then undergo the dark decay $p \rightarrow \tilde{\chi} e^+ \nu_e$. The mass

range $937.761 \text{ MeV} < m_\chi < 937.993 \text{ MeV}$ is problematic as well – although the proton remains stable, the ${}^9\text{Be}$ nucleus, known to be stable, can then undergo the nuclear dark decay ${}^9\text{Be} \rightarrow \chi + 2\alpha$ [30].

The mass range $937.993 \text{ MeV} < m_\chi < 938.783 \text{ MeV}$ is of particular interest, since then the proton, all stable nuclei and the dark matter χ remain stable, but the neutron can undergo the dark decay $n \rightarrow \bar{\chi} \gamma$. This decay channel was proposed in [31] as a possible solution to the neutron lifetime discrepancy. It was shown that Model 1B allows the neutron to have a dark decay branching fraction $\text{Br}(n \rightarrow \bar{\chi} \gamma) = 1\%$, which corresponds to $m_\phi \sim \mathcal{O}(100 \text{ TeV})$ for order one couplings; see [32] for a detailed review of this proposal along with its experimental signatures. It was recently argued that within the framework of Model 1A, the neutron dark decay branching fraction is constrained to be $\text{Br}(n \rightarrow \bar{\chi} \gamma) < 10^{-6}$ [33].

In Models 1A–1C, the dark matter χ can annihilate with the neutron, leading to signatures such as $\chi n \rightarrow \gamma + \text{meson(s)}$ at Super-Kamiokande and the future DUNE [34]. We will discuss this in more detail in the context of Model 2A below. Finally, in the case $m_\chi > m_p + m_e = 938.783 \text{ MeV}$ the dark particle χ is unstable, since the decay channel $\chi \rightarrow \bar{p} e^+ \nu_e$ opens up kinematically. Although in this scenario χ is not the dark matter, it can still be produced in experiments and lead to detectable signatures in B factories, e.g., missing energy signals from dark decays of heavy baryons and mesons [35].

Model 2A ($qqq\chi\phi$)

The Lagrangian for Model 2A is given by a simple extension of the interactions in Eq. (1),

$$-\mathcal{L}_2 \supset \lambda_q^{ab} \epsilon^{ijk} \Phi_i d_{Rj}^a d_{Rk}^b + \lambda_\chi^a \Phi^{*i} \tilde{\chi} u_{Ri}^a + \lambda_\phi \tilde{\chi} \chi \phi + \text{h.c.} \quad (2)$$

Once again, the coupling λ_q^{ab} is antisymmetric in flavor. The stability of the proton and all stable nuclei is guaranteed by the condition $m_\chi + m_\phi > 937.993 \text{ MeV}$. In the special case $937.993 \text{ MeV} < m_\chi + m_\phi < m_n = 939.565 \text{ MeV}$, the neutron can undergo the decay $n \rightarrow \bar{\chi} \phi^*$. This is the second neutron dark decay channel proposed in [31] to solve the neutron lifetime discrepancy, and leads to unique signals in nuclear decays [30]. In contrast to Model 1A, a large mass of χ or ϕ does not necessarily lead to dark matter decay; provided that $|m_\chi - m_\phi| < m_p + m_e = 938.783 \text{ MeV}$, both χ and ϕ remain stable.

Upon including an additional heavy particle $\tilde{\chi}'$ with the coupling $\tilde{\chi}' \chi \phi$, non-trivial CP phases can lead to different decay probabilities for $\tilde{\chi} \rightarrow udd$ versus $\tilde{\chi} \rightarrow \bar{u}d\bar{d}$ through the interference between tree-level and loop-level decay channels. Model 2A is then capable of explaining the matter-antimatter asymmetry of the universe through the hylogenesis mechanism [36] if $m_\chi + m_\phi \sim 5 \text{ GeV}$.

This model also predicts striking signatures in direct detection experiments: $\chi N \rightarrow \phi^* + \text{meson(s)}$, where N is a proton or a neutron. Such signals were studied in [37–39,34] and include: $\chi n \rightarrow \phi^* \pi^0$, $\chi p \rightarrow \phi^* \pi^+$, $\chi n \rightarrow \phi^* K^0$ and $\chi p \rightarrow \phi^* K^+$. Interestingly, not all of them are present in Model 2A. Given the antisymmetric structure of λ_q^{ab} in Eq. (2), the scalar mediator does not couple to two down quarks, which implies that signatures involving solely pions in the final state are not possible in Model 2A.

In conclusion, we find that there is a specific prediction of baryonic dark matter models without tree-level proton decay to Standard Model particles: dark matter-nucleon annihilation leads to at least one kaon in the final state. This is an especially relevant observation in light of DUNE's expected exquisite sensitivity to kaons. In certain regions of parameter space, DUNE will be able to probe Φ masses up to $m_\phi \sim \mathcal{O}(100 \text{ TeV})$ [34].

Table 2

Effective operators describing the interaction of leptonic dark matter with leptons and their model realizations.

Model	Leptonic dark matter	
	Interactions	Mediators
$\bar{l}l\chi$		
3A	$\Phi L_L L_L, \Phi^* \bar{e}_R \chi$	$\Phi = (1, 1)_1$
3B	$\Phi L_L \bar{e}_R, \Phi^* L_L \chi$	$\Phi = (1, 2)_{-\frac{1}{2}}$
$\bar{l}l\chi\phi$		
4A	$\Phi L_L L_L, \Phi^* \bar{e}_R \tilde{\chi}, \tilde{\chi} \chi \phi$	$\Phi = (1, 1)_1 \quad \tilde{\chi} = (1, 1)_0$
4B	$\Phi L_L L_L, \Phi^* \tilde{\chi} \chi, \tilde{\chi} \bar{e}_R \phi$	$\Phi = (1, 1)_1 \quad \tilde{\chi} = (1, 1)_1$
4C	$\Phi L_L \bar{e}_R, \Phi^* L_L \tilde{\chi}, \tilde{\chi} \chi \phi$	$\Phi = (1, 2)_{-\frac{1}{2}} \quad \tilde{\chi} = (1, 1)_0$
4D	$\Phi L_L \bar{e}_R, \Phi^* \tilde{\chi} \chi, \tilde{\chi} L_L \phi$	$\Phi = (1, 2)_{-\frac{1}{2}} \quad \tilde{\chi} = (1, 2)_{-\frac{1}{2}}$

Baryoleptonic dark matter

The simplest class of baryoleptonic dark matter models arises from dimension seven effective operators $qqql\phi$ and $qqq\bar{l}\phi$, where l is a Standard Model lepton representation. Minimal particle physics realizations of $qqql\phi$ involve one of the scalar mediators $(3, 1)_{-1/3}$, $(3, 1)_{-4/3}$, or one of the vectors $(3, 2)_{-5/6}$, $(3, 2)_{1/6}$, whereas particle models for $qqq\bar{l}\phi$ require the scalar mediator $(3, 1)_{-1/3}$ or the vector $(3, 2)_{1/6}$. In models with only a single mediator, it has to couple to the bilinears qq and ql (or $q\bar{l}$), resulting in tree-level proton decay which cannot be forbidden by any symmetry.

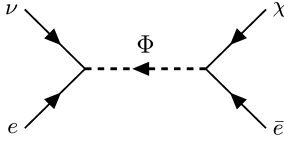
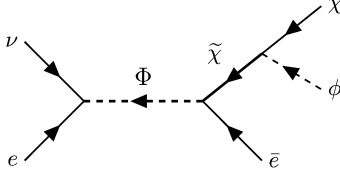
A possible way to overcome this issue is to introduce two mediators, one coupled only to qq , and the other coupled only to ql (or $q\bar{l}$) [35], or include an additional heavy fermion [40]. However, we will not consider those scenarios here, since they require imposing baryon/lepton number conservation. A simple extension of this class of models, described by the effective operator $qqql\phi^2$, i.e., replacing ϕ by ϕ_e and adding the interaction term $\phi_e^* \phi^2$ to stabilize ϕ , provides a working mechanism for baryogenesis as in Model 2A and also exhibits nonstandard nucleon destruction signatures [40]. A more general analysis of the baryon asymmetry generation in those types of models was performed in [41].

Leptonic dark matter

In order to explain the matter-antimatter asymmetry through leptogenesis, operators carrying a nonzero lepton number contribution from the leptons are needed. Those operators can generate an asymmetry in the lepton sector, which is then transferred to the baryon sector by electroweak sphalerons [42]. The dimension four and five operators $\bar{l}l\phi$ and $\bar{l}l\phi^2$ are not of interest to us since ϕ does not carry lepton number. Also, as in the baryonic dark matter case, we will not consider operators involving the Higgs, e.g., the dimension four operator $HL_L\chi$.

The first operator of interest is the dimension six $\bar{l}l\chi$, where l can be either L_L or e_R . The only gauge-invariant realization is $L_L L_L \bar{e}_R \chi$. Note that the operator $\bar{l}l\chi$ is not invariant under hypercharge. There are two particle models for the operator $\bar{l}l\chi$, denoted as Models 3A and 3B in Table 2. The corresponding mediators are the scalars $(1, 1)_1$ and $(1, 2)_{-1/2}$, respectively. They are colorless, thus they do not mediate proton decay. Because of the similarity between the two models, we write down the Lagrangian only for Model 3A. Its realization of the operator $\bar{l}l\chi$ is shown in Fig. 3.

At dimension seven, the effective operator of interest is $\bar{l}l\chi\phi$, again with only a single gauge-invariant realization $L_L L_L \bar{e}_R \chi \phi$. The corresponding particle models are labeled as Models 4A–4D in Table 2. An intermediate particle $\tilde{\chi}$ is required. We focus on Model 4A below for a quantitative discussion of its properties. The re-

Fig. 3. Model 3A realization of the operator $ll\bar{l}\chi$.Fig. 4. Model 4A realization of the operator $ll\bar{l}\chi\phi$.

alization of the operator $ll\bar{l}\chi\phi$ within this model is presented in Fig. 4.

Regarding other options, not involving the Higgs field, the dimension seven operator $llll\phi$ does not carry nonzero lepton number, thus it cannot fit into the framework of asymmetric dark matter. At dimension eight, the only operator carrying lepton charge is $ll\bar{l}\chi\phi^2$. As in the baryonic dark matter case, models for this operator are obtained from Models 4A–4D by substituting ϕ with ϕ_e , and introducing the interaction $\phi_e^*\phi^2$. This makes ϕ automatically stable, but offers no other advantages compared to the $ll\bar{l}\chi\phi$ case.

Model 3A ($ll\bar{l}\chi$)

The Lagrangian for Model 3A is

$$-\mathcal{L}_3 \supset \lambda_l^{ab} \Phi (L_L^a \epsilon L_L^b) + \lambda_\chi^a \Phi^* \chi \bar{e}_R^a + \text{h.c.}, \quad (3)$$

where again a, b are flavor indices and the parenthesis denotes the contraction of $SU(2)_L$ indices. Because of the antisymmetric structure of the LL bilinear, the coupling λ_l^{ab} must be antisymmetric in flavor.

This model is not phenomenologically attractive because it is hard to ensure the stability of the dark matter particle. Especially, without imposing any lepton flavor symmetry, χ can undergo the decay $\chi \rightarrow \nu\gamma$ through a loop diagram. Thus, we do not consider this class of models further.

Model 4A ($ll\bar{l}\chi\phi$)

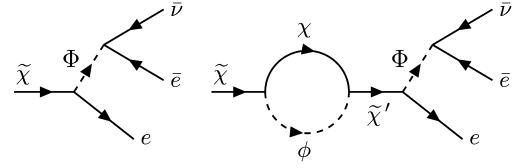
A phenomenologically viable model which contains a dark matter candidate and can successfully explain the matter-antimatter of the universe is obtained by introducing an additional scalar particle ϕ . The resulting Lagrangian is a minimal extension of \mathcal{L}_3 ,

$$-\mathcal{L}_4 \supset \lambda_l^{ab} \Phi (L_L^a \epsilon L_L^b) + \lambda_\chi^a \Phi^* \chi \bar{e}_R^a + \lambda_\phi \bar{\chi} \chi \phi + \text{h.c.} \quad (4)$$

The same symmetry arguments as for Model 3A apply, i.e., λ_l^{ab} is antisymmetric in flavor. In addition, we imposed a Z_2 parity in the dark sector, under which ϕ and χ are odd, so that the lighter of them remains stable and can constitute the dark matter. Model 4A is the subject of the subsequent section.

3. Phenomenology of leptonic dark matter

The crucial property of leptonic dark matter models is that they do not suffer from proton decay. Below we discuss the baryon asymmetry generation in Model 4A via leptogenesis, the symmetry restoration in the dark sector and the signatures expected in indirect detection experiments. Our conclusions apply to Models 4B–4D as well.

Fig. 5. Diagrams contributing to the decay $\tilde{\chi} \rightarrow e_R \bar{e}_L \bar{\nu}_L$.

Leptogenesis

The generation of lepton asymmetry within the framework of Model 4A is similar to the baryogenesis mechanism [36]. The process starts immediately after inflation. Once the inflaton field Ψ falls into the potential minimum, it starts oscillating and leads to the reheating of the universe. We assume that during reheating the particles $\tilde{\chi}$ and $\bar{\tilde{\chi}}$ are produced in equal amounts. As the temperature drops, the particles $\tilde{\chi}$ decay via the channels $\tilde{\chi} \rightarrow \chi\phi$ and $\tilde{\chi} \rightarrow e_R \bar{e}_L \bar{\nu}_L$, where \bar{e}_L and $\bar{\nu}_L$ have different flavors. The antiparticles $\bar{\tilde{\chi}}$ decay through conjugate processes.

In order to generate CP violation in the model, we introduce the particle $\tilde{\chi}'$ with the following interaction terms,

$$-\mathcal{L}'_4 \supset \lambda_{\tilde{\chi}'}^a \Phi^* \tilde{\chi}' \bar{e}_R^a + \lambda_{\phi}^{\tilde{\chi}'} \bar{\tilde{\chi}'} \chi \phi + \text{h.c.}, \quad (5)$$

and we take $m_{\tilde{\chi}'} \gg m_{\tilde{\chi}}$. The asymmetry between the decays of $\tilde{\chi}$ and $\bar{\tilde{\chi}}$ arises from the interference between the tree-level and one-loop diagram shown in Fig. 5. If the leading decay channel is $\tilde{\chi} \rightarrow \chi\phi$, the generated lepton asymmetry is¹

$$\begin{aligned} \Delta_L &= \frac{\Gamma(\tilde{\chi} \rightarrow e_R \bar{e}_L \bar{\nu}_L) - \Gamma(\bar{\tilde{\chi}} \rightarrow \bar{e}_R e_L \nu_L)}{2 \Gamma(\tilde{\chi} \rightarrow \chi \phi)} \\ &\approx \frac{|\lambda_l|^2 \text{Im}(\lambda_{\tilde{\chi}'}^* \lambda_{\tilde{\chi}} \lambda_{\phi} \lambda_{\phi}^*)}{1536 \pi^3 |\lambda_{\phi}|^2} \frac{m_{\tilde{\chi}}^5}{m_{\tilde{\chi}'}^4 m_{\tilde{\chi}}}. \end{aligned} \quad (6)$$

In order to avoid the washout of the asymmetry by $\chi\phi \rightarrow e_R \bar{e}_L \bar{\nu}_L$ scattering, the reheating temperature, for couplings $\mathcal{O}(1)$, needs to satisfy [36]

$$T_R \lesssim (11 \text{ TeV}) \left[\frac{m_{\phi}^4 m_{\tilde{\chi}}^2}{(1 \text{ PeV})^6} \right]^{\frac{1}{5}}. \quad (7)$$

Provided that T_R is above the electroweak symmetry breaking scale, i.e., $T_R \gtrsim 200 \text{ GeV}$, a portion of the created lepton asymmetry will be converted into a baryon asymmetry by the electroweak sphalerons. Therefore, as long as the masses of Φ and $\tilde{\chi}$ fulfill the condition

$$(m_{\phi}^4 m_{\tilde{\chi}}^2)^{\frac{1}{6}} \gtrsim 40 \text{ TeV}, \quad (8)$$

the sphalerons efficiently transfer the asymmetry to the baryonic sector. We find that the final baryon asymmetry is

$$\Delta_B \approx \frac{28}{79} \Delta_L. \quad (9)$$

The exact relation between the baryon-to-photon ratio η_B and the baryon asymmetry depends on the model of reheating; up to $\mathcal{O}(1)$ factors it is given by [43–45]

$$\eta_B \approx \frac{\Delta_B T_R}{M_{\Psi}}, \quad (10)$$

where M_{Ψ} is the inflaton mass. The observed value of the baryon-to-photon ratio in the universe of $\eta_B \approx 6 \times 10^{-10}$ [46] is obtained,

¹ In this calculation, we only consider the decay channel with one particular lepton flavor choice. The rescaling can be easily done to include more flavor channels.

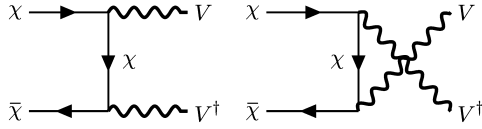


Fig. 6. Annihilation of the dark matter symmetric component.

e.g., for $\mathcal{O}(1)$ couplings, $m_\phi = m_{\tilde{\chi}} = 1$ PeV, $m_{\tilde{\chi}'} = 5$ PeV, $M_\Psi = 20$ PeV and $T_R = 8$ TeV.

The sum of χ and ϕ masses is set by the observed ratio of the dark matter and baryon abundances. Since the generated dark matter asymmetry $\Delta_{\text{DM}} = \Delta_L$, one arrives at

$$m_\chi + m_\phi = m_p \frac{\Omega_{\text{DM}}}{\Omega_B} \left| \frac{\Delta_B}{\Delta_{\text{DM}}} \right| \approx 1.8 \text{ GeV}. \quad (11)$$

Removing the symmetric component

To ensure the annihilation of the symmetric component of dark matter, it is sufficient to introduce one new particle, e.g., a vector V , lighter than χ and ϕ , with the interaction terms

$$-\mathcal{L}'_4 \supset g_\chi V_\mu \bar{\chi} \gamma^\mu \chi + g_\phi V_\mu \phi \partial^\mu \phi^* + g_V V_\mu \bar{e}_R^a \gamma^\mu e_R^a + \text{h.c.} \quad (12)$$

This allows for the annihilation channels $\bar{\chi}\chi \rightarrow VV^\dagger$ and $\phi\phi^* \rightarrow VV^\dagger$, along with the subsequent decays $V \rightarrow l^+l^-$. Diagrams corresponding to $\bar{\chi}\chi \rightarrow VV^\dagger$ are shown in Fig. 6. We note that the cross section for the annihilation $\bar{\chi}\chi \rightarrow l^+l^-$ through an s -channel V is suppressed due to stringent experimental constraints on the coupling g_V [47].

The annihilation cross section $\sigma_{\bar{\chi}\chi}$ is related to the symmetric component of the χ relic density as

$$\Omega_\chi h^2 \simeq \left(\frac{9 \times 10^{-11}}{\text{GeV}^2} \right) \frac{1}{\langle \sigma_{\bar{\chi}\chi} v \rangle \sqrt{g_*}} \frac{m_\chi}{T_f}, \quad (13)$$

where g_* is the number of relativistic degrees of freedom and T_f is the freeze-out temperature [48]. Given the observed value of the dark matter relic density $\Omega_{\text{DM}} h^2 = 0.12$ [49], we find that for $m_V \ll m_\chi$ an efficient annihilation of the symmetric component of χ is achieved when

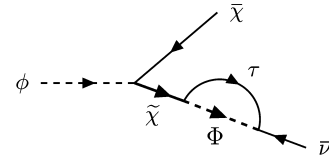
$$g_\chi \gtrsim 0.01. \quad (14)$$

Similar arguments apply to the annihilation $\phi\phi^* \rightarrow VV^\dagger$. The decays of V , such as $V \rightarrow e^+e^-$, depending on its coupling and mass, can be a slow process. Thus a small coupling between V and leptons is allowed, and can easily be consistent with all experimental constraints [47].

Another possibility is to introduce a light scalar ϕ' , rather than the vector boson V , to remove the symmetric components of χ and ϕ . The annihilation $\bar{\chi}\chi \rightarrow \phi'\phi'^*$ is a p -wave process, thus it is suppressed by the relative velocity. Such a velocity suppression does not significantly affect the annihilation of the symmetric components, since both χ and ϕ are still semi-relativistic during their freeze-out. However, it has interesting implications for the experimental constraints and signatures, which will be discussed below.

Symmetry restoration

Interestingly, within Model 4A the asymmetry in the dark sector is not preserved during the evolution of the universe. If $m_\phi > m_\chi$, the interactions in Eq. (4) render the particle ϕ unstable, resulting in the two-body decay $\phi \rightarrow \bar{\chi}\bar{\nu}$ and the four-body decay

Fig. 7. Diagram for the decay $\phi \rightarrow \bar{\chi}\bar{\nu}$.

$\phi \rightarrow \bar{\chi}\bar{\nu}e^+e^-$. The two-body channel is dominant and proceeds through the diagram shown in Fig. 7. The resulting decay rate is

$$\Gamma(\phi \rightarrow \bar{\chi}\bar{\nu}) \sim \frac{|\lambda_l \lambda_{\tilde{\chi}} \lambda_\phi|^2}{4096 \pi^5} \frac{m_\tau^2 m_\phi}{m_{\tilde{\chi}}^2} \left(1 - \frac{m_\chi^2}{m_\phi^2} \right)^2. \quad (15)$$

The observed dark matter relic density imposes a constraint on this rate. More explicitly, the decays $\phi \rightarrow \bar{\chi}\bar{\nu}$ restore the symmetry between χ and $\bar{\chi}$ in the relic abundance. Such a restoration must happen sufficiently late so that χ and $\bar{\chi}$ do not efficiently annihilate with each other to cause $\mathcal{O}(1)$ change to the dark matter relic abundance. We find that this requirement is met if the decay happens at temperatures

$$T_{\phi \text{ decay}} \lesssim 50 \text{ MeV} \quad (16)$$

or, equivalently, after $t_\phi \simeq \mathcal{O}(10^{-4})$ s. This can be achieved if the mass splitting between χ and ϕ is small, i.e., $m_\phi \approx m_\chi \approx 0.9$ GeV. For example, taking $\mathcal{O}(1)$ couplings, $m_{\tilde{\chi}} \sim 1$ PeV and $(m_\phi - m_\chi)/m_\phi \sim 10^{-3}$, the lifetime of ϕ is $\tau_\phi \sim 0.1$ s. Since the decays of ϕ 's only produce slow-moving χ 's and low-energy neutrinos due to the small mass splitting, the cosmological constraints, e.g., from Big Bang Nucleosynthesis, are easily evaded [50,51].

The restored symmetry in the dark sector revives $\bar{\chi}\chi$ annihilation at present times in regions with large dark matter concentration, e.g., in the Galactic Center, which does not happen in standard asymmetric dark matter models. A similar scenario of symmetry restoration in the dark sector was proposed in the context of heavy dark matter [18–20] and oscillating dark matter [21–25]. In addition, late decays of the heavier dark matter component, $\tau_\phi \sim 10^9$ years, along with a small mass splitting between the components, are consistent with observation [52,53]. Such a late decay is also proposed as a possible solution to the missing satellites problem [54] and the core-cusp problem [55].

Indirect detection signatures

The symmetry restoration in the dark sector induces nontrivial dark matter indirect detection signals. For example, if the mediator V couples to leptons and its mass $m_V > 2m_\mu$, the dark matter annihilation leads to final states involving four leptons, i.e., $e^+e^-e^+e^-$, $e^+e^-\mu^+\mu^-$ or $\mu^+\mu^-\mu^+\mu^-$. Those leptons can further undergo inverse Compton scattering and bremsstrahlung, producing photons that can be measured in indirect detection experiments.

The present-day cross section for the dark matter annihilation via $\bar{\chi}\chi \rightarrow VV^\dagger \rightarrow l^+l^-l^+l^-$ expected in Model 4A is $\langle \sigma_{\bar{\chi}\chi} v \rangle \gtrsim 10^{-26} \text{ cm}^3/\text{s}$ and corresponds to the red region in Fig. 8. The minimum value for the coupling g_χ in Eq. (12) leading to an efficient annihilation of the symmetric component of χ in the early universe is ~ 0.01 (see Eq. (14)), which results in the cross section $\langle \sigma_{\bar{\chi}\chi} v \rangle \approx 10^{-26} \text{ cm}^3/\text{s}$ in the present epoch, denoted by the black dot.

The most stringent constraint on $\langle \sigma_{\bar{\chi}\chi} v \rangle$ arises from the measurements of the Cosmic Microwave Background provided by the Planck satellite. The resulting bound [56] is shown as the brown curve in Fig. 8. For $m_\chi \approx 0.9$ GeV, which is required in Model 4A

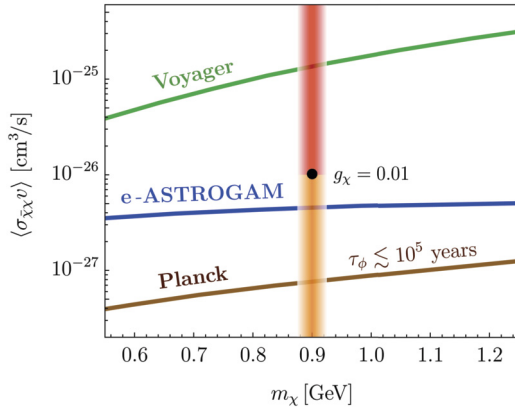


Fig. 8. Dark matter annihilation cross section in Model 4A expected in the present epoch when the annihilation is mediated by the vector V (red region) or by the scalar ϕ' (orange and red regions). Overplotted are the exclusion limits: the green curve corresponds to the bound on $\bar{\chi}\chi \rightarrow e^+e^-$ from Voyager [59], whereas the brown curve corresponds to the bound on $\bar{\chi}\chi \rightarrow e^+e^-e^+e^-$ from Planck [56] applicable only when $\tau_\phi \lesssim 10^5$ years. The projected future e-ASTROGAM reach for $\bar{\chi}\chi \rightarrow e^+e^-e^+e^-$ [61] is denoted by the blue curve. The black dot corresponds to the benchmark scenario $g_\chi = 0.01$.

due to Eqs. (11) and (16), the annihilation cross section needs to be $\langle\sigma_{\bar{\chi}\chi}v\rangle \lesssim 7 \times 10^{-28} \text{ cm}^3/\text{s}$. However, this constraint can be evaded in two ways: (1) if ϕ or the light mediator V is very long-lived due to either a small mass splitting or tiny couplings, in which case the energy deposition to the Standard Model thermal bath only happens after the recombination, or (2) if the light mediator is a scalar ϕ' , since the dark matter annihilation cross section is then largely reduced due to velocity suppression at late times.

In scenario (1), i.e., if $\tau_\phi \gg 10^5$ years, the Planck bound does not apply and the most stringent constraint on the annihilation cross section for $m_\chi \approx 0.9$ GeV comes from the data collected by Voyager [57,58]. The green curve in Fig. 8 shows the Voyager bound assuming that the dominant annihilation channel is $\bar{\chi}\chi \rightarrow e^+e^-$ [59]. The bound on the cross section for $\bar{\chi}\chi \rightarrow l^+l^-l^+l^-$ (for $m_\chi \approx 0.9$ GeV) has not been determined, but it is expected to be weaker, as demonstrated for heavier dark matter based on the Fermi Gamma-Ray Space Telescope data [60]. Therefore, most of the parameter space region shown in red in Fig. 8 remains phenomenologically viable.

In scenario (2) the Planck bound does apply, however, due to velocity suppression, the dark matter annihilation cross section can be substantially smaller at present times compared to case (1), as indicated by the orange region in Fig. 8. The cross section can be as small as $\langle\sigma_{\bar{\chi}\chi}v\rangle \sim 10^{-32} \text{ cm}^3/\text{s}$. Although the parameter space above the brown line is excluded in this scenario, there still remains a vast region below the Planck bound which has not been explored.

Interestingly, the entire parameter space for scenario (1), corresponding to the red region in Fig. 8, will be probed by the future e-ASTROGAM experiment, whose predicted reach is shown as the blue curve. For $m_\chi \approx 0.9$ GeV, e-ASTROGAM is expected to be sensitive to dark matter annihilation cross sections $\langle\sigma_{\bar{\chi}\chi}v\rangle \gtrsim 5 \times 10^{-27} \text{ cm}^3/\text{s}$ [61].

4. Summary

The nature of dark matter, origin of the matter-antimatter asymmetry of the universe and proton stability are certainly among the greatest open questions in modern particle physics. Theories in which dark matter couples to quarks and leptons introduce a natural framework for solving the first two of those puzzles. However, many of such models suffer from tree-level proton decay

to Standard Model particles unless a conservation of baryon and/or lepton number is imposed by hand.

In this letter, we chose proton stability as the primary criterion for a model's naturalness. We demonstrated that this reduces the number of viable baryonic dark matter theories to just a few models involving the scalar mediator $(3, 1)_{2/3}$ and a sufficiently heavy dark matter particle. This has an interesting impact on the potential experimental signatures – we found that for baryonic models naturally free from proton decay, the final state of dark matter-nucleon annihilation necessarily involves a kaon.

We also considered a model of leptonic dark matter, which explains the matter-antimatter asymmetry of the universe through leptogenesis and predicts symmetry restoration in the dark sector. In this theory the dark matter annihilation may be enhanced at late times, providing signals that can be searched for in future indirect detection experiments.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

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