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Secluded WIMP dark matter

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Abstract

We consider a generic mechanism via which thermal relic WIMP dark matter may be decoupled from the Standard Model, namely through a combination of WIMP annihilation to metastable mediators with subsequent delayed decay to Standard Model states. We illustrate this with explicit examples of WIMPs connected to the Standard Model by metastable bosons or fermions. In all models, provided the WIMP mass is greater than that of the mediator, it can be secluded from the Standard Model with an extremely small elastic scattering cross-section on nuclei and rate for direct collider production. In contrast, indirect signatures from WIMP annihilation are consistent with a weak scale cross-section and provide potentially observable γ-ray signals. We also point out that γ-ray constraints and flavor physics impose severe restrictions on MeV-scale variants of secluded models, and identify limited classes that pass all the observational constraints.

1. Introduction

The overwhelming astrophysical and cosmological evidence for dark matter has in recent years led to a dramatic expansion in experimental programs that aim to detect observational signatures of its non-gravitational interactions [1]. These probes range from direct production at colliders to the recoil of galactic dark matter on nuclei in underground detectors, and indirect detection of annihilation products, primarily γ-rays. A driving paradigm in developing these probes is that of WIMP (weakly interacting massive particle) dark matter, which represents a simple and attractive candidate through the fact that a thermal relic with weak scale mass and annihilation cross-section into Standard Model (SM) states naturally provides roughly the correct cosmological abundance [2]. This weak-scale annihilation cross-section, when reversed, naturally suggests a weak-scale production cross-section at colliders, and when viewed in the t-channel implies an elastic scattering cross-section on nuclei which is within reach of purpose-built underground detectors. In recent years, experiments of the latter type have reached an impressive level of sensitivity [3,4], and significant future progress in this direction is anticipated.

Along with direct scattering on nuclei, and the neutrino signal from the annihilation of WIMPs inside the solar core, which is sensitive to the WIMP trapping rate (and again to the elastic scattering cross-section), the indirect detection of WIMPs via the products of their annihilation in the center of the galaxy is a distinct possibility [5]. This annihilation signal can be very model dependent, due to uncertainties in the halo profile, varying branching ratios, and the presence or absence of detectable monoenergetic photons. Thus, prior to specifying a particular model, it is impossible to say which strategy, direct or indirect detection, will provide a more sensitive probe of WIMPs. However, within the prevailing WIMP paradigm as exemplified by the lightest superpartner (LSP) in the MSSM for example, the stringent constraints on the elastic scattering cross-section often impose significant limits on the available indirect signal from annihilation.

In this Letter, we will revisit this aspect of WIMP physics, and question the commonplace assumption of a close link be-

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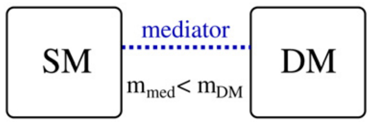


Fig. 1. The secluded WIMP dark matter scenario.

tween the cross-sections relevant for direct and indirect detection. More precisely, we consider what constraints the picobarn annihilation cross-section required of a thermal relic WIMP can actually impose on its interactions with normal matter, e.g. production at colliders or scattering off nuclei. We observe that in relatively simple models, the latter interactions can be highly suppressed and thus in many cases the indirect annihilation signature will be the most important probe of the non-gravitational interactions of dark matter.

Generically, any WIMP dark matter model can be conveniently decomposed into three sectors, the SM, the WIMP itself, and the fields which mediate the WIMPs interactions with the SM,

$$\mathcal{L} = \mathcal{L}_{\mathrm{SM}} + \mathcal{L}_{\mathrm{WIMP}} + \mathcal{L}_{\mathrm{mediator}}, \tag{1}$$

and in many models the mediator states are in fact part of the SM, e.g. the electroweak gauge bosons or the Higgs. In this Letter, we point out a simple and generic mechanism that allows the WIMP, while still a thermal relic, to be *secluded* from the SM, dramatically reducing its couplings to SM states, and consequently suppressing the collider and direct detection rates by many orders of magnitude. The mechanism relies on a metastable mediator that couples the SM to the secluded WIMP sector, with a mass less than that of the WIMP, see Fig. 1. In this kinematic regime, direct annihilation into a pair of mediators is always possible. Due to their coupling to the Standard Model, these particles are unstable but the constraint on their lifetime is very weak, and in particular a lifetime under one second is sufficient to guarantee their decay before the beginning of primordial nucleosynthesis (BBN), rendering them completely harmless. Secluded WIMP models may therefore be impossible to detect using colliders or direct searches, but the indirect signatures, through e.g. annihilation in the Galactic center, can be as pronounced as in any WIMP scenario.

To illustrate this mechanism, in Section 2 we will construct several models with fermionic, scalar and vector particles as mediators. We show that if $m_{\mathrm{WIMP}} < m_{\mathrm{mediator}}$, the parameter space of such models is highly constrained, as the coupling of the mediator to the SM must necessarily be sizable to ensure the required annihilation cross-section. Yet if the reverse is true, $m_{\mathrm{WIMP}} > m_{\mathrm{mediator}}$, there are no strict requirements on the size of the mixing except for the lifetime of the mediator state, which in some instances can be satisfied for (mixing)^2 of the mediator with the SM as low as 10^{-23}.

An interesting limiting regime of the secluded scenario arises when the mediator (and the WIMP) are both very light. This ties in with models of MeV-scale dark matter that provide a tantalizing yet speculative link [6] to the unexpectedly strong 511 keV line observed from the galactic center [7]. We show in Section 3 that some secluded WIMPs may have advantages over existing models of this type, and discuss their observational signatures at low energies, including missing energy signals in meson decays.

2. Models of secluded WIMPs

For some time the discussion of WIMPs has centered around supersymmetric models where the LSP is stable, provided that R-parity is unbroken. Supersymmetry is largely motivated by a well-known combination of theoretical arguments unrelated to dark matter, and the possible existence of a stable or long-lived WIMP-LSP may provide an interesting bridge to cosmology. On the other hand, to date there are no experimental indications of supersymmetry, while there is ample evidence for the existence of dark matter. Therefore, an alternative approach to the particle physics of dark matter that is certainly logical and justifiable, consists of studying generic classes of WIMP models, among which the minimal choices are obviously well motivated. Over the years, WIMP models with Higgs and/or singlet mediation have been studied extensively [8–10]. More recently, models with exotic electroweak matter were also considered in some detail [12], as well as models with additional gauge groups [13]. More generally, going beyond the WIMP framework also allows for freedom in tuning the coupling to the SM, as in scenarios with sterile neutrinos [11] or dark matter populated by late decays [14]. In this Letter we adhere to a rather minimalist WIMP framework, which is well suited to demonstrating our main point.

2.1. *U(1)′ mediator*

We can construct a simple secluded model, starting from the Lagrangian for a Dirac WIMP, whose interaction with the SM is mediated by an additional U(1)′ gauge group:

$$\begin{aligned}\mathcal{L}_{\mathrm{WIMP+mediator}} = &-\frac{1}{4}V_{\mu\nu}^2 - \frac{\kappa}{2}V_{\mu\nu}B_{\mu\nu} - |D_\mu\phi|^2 - U\left(\phi\phi^*\right)\\ &+ \bar{\psi}(iD_\mu\gamma_\mu - m_\psi)\psi.\end{aligned} \tag{2}$$

In this Lagrangian, ψ and V_μ denote respectively the Dirac WIMP and the U(1)′ vector boson mediator, with field strength $V_{\mu\nu}$ and covariant derivative $D_\mu = \partial_\mu + ie'V_\mu$. To avoid confusion, we denote the strength of the U(1)′ coupling constant as e'. The U(1)′ vector bosons V_μ couple to the SM hypercharge gauge bosons B_μ via kinetic mixing. The additional scalar ϕ higgses U(1)′ at or near the weak scale. Note that the SM and WIMP sectors are coupled only via the mediator at the renormalizable level, the SM is neutral under U(1)′, and the WIMP sector is a singlet under the SM gauge group. Assuming the scalar potential U has a minimum away from zero, the Lagrangian below the U(1)′ breaking scale takes the form,

$$\begin{aligned}\mathcal{L}_{\mathrm{WIMP+mediator}} = &-\frac{1}{4}V_{\mu\nu}^2 + \frac{1}{2}m_V^2V_\mu^2 + \kappa V_\nu\partial_\mu B_{\mu\nu}\\ &+ \bar{\psi}(iD_\mu\gamma_\mu - m_\psi)\psi + \mathcal{L}_{h'},\end{aligned} \tag{3}$$

where m_V is the resulting mass of the U(1)′ vector boson, and $\mathcal{L}_{h'}$ is the Lagrangian for the Higgs′ particles including their self-interaction and interactions with V_μ.

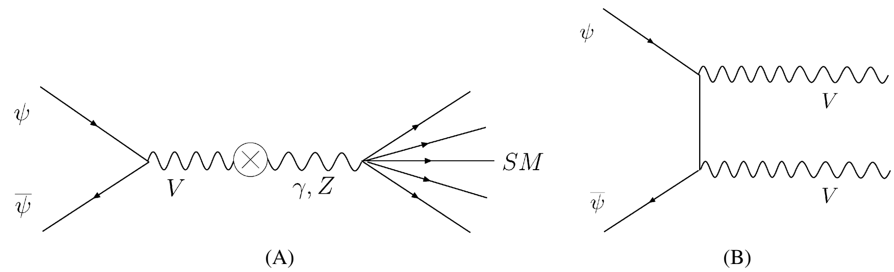


Fig. 2. WIMP annihilation for: (A) $m_\psi < m_V$ on the left; and (B) $m_\psi > m_V$ on the right—the secluded regime in which the annihilation may proceed via two metastable on-shell V's, which ultimately decay to SM states.

There are four important parameters in the model, the WIMP and vector boson masses m_ψ and m_V, the mixing parameter κ and new coupling constant e'. The most important quantity for WIMP physics is arguably the annihilation cross-section into the SM states. To this end it is easy to identify the two primary mechanisms responsible for annihilation (see Fig. 2):

(A) $\psi + \bar{\psi} \to$ virtual $V \to$ virtual $\gamma, Z \to$ SM states.
(B) $\psi + \bar{\psi} \to$ on-shell $V + V$, with subsequent decay to SM states.

Process (B) is open only in the kinematic regime $m_\psi > m_V$ while process (A) can occur regardless of the relation between m_V and m_ψ. We will analyze the case of $m_\psi < m_V$ first.

2.1.1. *The characteristic WIMP regime*

The annihilation cross-section in this case is given by the diagram in Fig. 2(A). Although we use the full result for numerical analysis, to simplify the presentation its helpful to quote the annihilation cross-section in the limit $m_Z^2, m_t^2, m_h^2 \ll m_\psi^2$. Since $2m_\psi$ then provides the energy scale for the problem, in this limit one may substitute $\partial_\mu B_{\mu\nu}$ by the total hypercharge current and neglect the influence of SM threshold effects. For small mixing, characterized by $\beta \ll 1$ where

$$\beta \equiv \left(\frac{\kappa e'}{e\cos\theta_W}\right)^2, \tag{4}$$

the resulting annihilation cross-section for nonrelativistic WIMPs takes the following form,

$$\langle\sigma_{\rm ann} v\rangle_{m_\psi \gg m_{\rm SM}}$$
$$\approx 1.3\ {\rm pbn} \times \beta\left(\frac{500\ {\rm GeV}}{m_\psi}\right)^2 \times \left(\frac{4m_\psi^2}{4m_\psi^2 - m_V^2}\right)^2, \tag{5}$$

proceeding in the $l = 0$ channel with an obvious pole at $m_\psi = m_V/2$, in the vicinity of which a more accurate treatment of the thermal average is required. The result depends on the mixing parameter β and the sum of squares of the hypercharges for the SM fields, $\sum_{\rm fermions} Y_f^2 + \frac{1}{2}\sum_{\rm bosons} Y_b^2 = 10 + 0.25$. Note that in the opposite limit, $m_b \ll m_\psi \ll m_Z$, the total cross-section is instead proportional to the sum of squares of all the electric charges of SM fermions with the exception of the t-quark.

This cross-section needs to be compared with the constraint on the dark matter energy density provided by recent cosmological observations:

$$2 \times \frac{10^9 (m_\psi/T_f)}{\sqrt{g_*(T_f)} \times {\rm GeV} \times M_{\rm Pl}\langle\sigma v\rangle} \leqslant \Omega_{\rm DM} h^2 \simeq 0.1, \tag{6}$$

where T_f is the freeze-out temperature (it suffices here to take $m_\psi/T_f \simeq 20$), g_* the effective number of degrees of freedom at freeze-out, and the extra factor of two relative to the standard formula (see e.g. [16]) is because annihilation can occur only between particles and anti-particles.

In Fig. 3, we exhibit the abundance constraint on the β–m_ψ plane for a specific choice of mediator mass, $m_V = 400$ GeV, by saturating the inequality (6). This value of m_V lies outside the direct reach of LEP or the Tevatron but is certainly within range for the LHC. One can clearly see the enhancement of the annihilation cross-section in the vicinity of the two vector resonance poles, Z and V, where the mixing parameter β is allowed to be significantly smaller than 1.

This model is subject to various constraints from direct searches and collider physics.

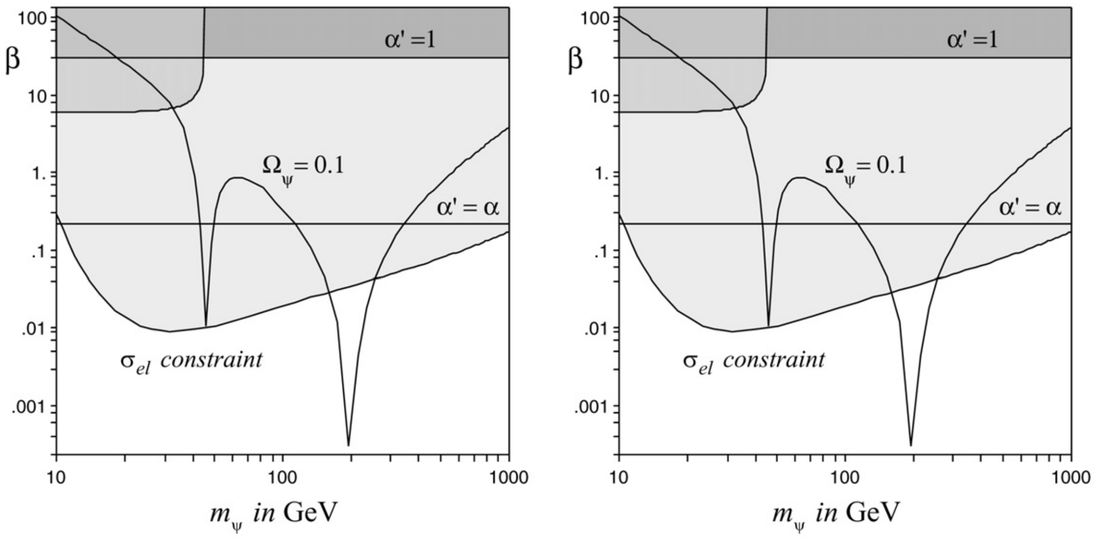
(a) The elastic scattering of galactic WIMPs off nuclei occurs with a characteristic momentum transfer of order 100 MeV or less, making virtual photon exchange the dominant contribution to the scattering amplitude. From an appropriate low-energy effective theory viewpoint, the WIMP ψ is electrically neutral but exhibits a non-zero charge radius given by

$$r_c^2 = 6\frac{\kappa e'}{e} \times \frac{1}{m_V^2} = 6\frac{\beta^{1/2}\cos\theta_W}{m_V^2}. \tag{7}$$

The contribution of r_c to the elastic scattering of WIMPs off nuclei was calculated in [17], and can easily be rescaled to the equivalent cross-section per nucleon,

$$\sigma_{\rm el} = \frac{4\pi}{9} Z^2\alpha^2 r_c^4 \left(\frac{m_A m_D}{m_A + m_D}\right)^2$$
$$\Longrightarrow \quad \sigma_{\rm nucleon} = \frac{4\pi}{9}\alpha^2 r_c^4 m_p^2 \left(\frac{Z}{A}\right)^2, \tag{8}$$

where m_A and m_p are nuclear and nucleon masses, and the Z/A ratio should be specified for the relevant experimental setup. We plot the corresponding experimental limit recently obtained by the XENON Collaboration [4] in Fig. 3, which clearly rules out a significant portion of the parameter space away from the $m_\psi = M_V/2$ resonance. This is not surprising, as the model is in many ways similar to the original "heavy Dirac neutrino" of



(A) (B)

Fig. 3. (A) On the left, the parameter space of the U(1)$'$-mediated model. The elastic cross-section constraint excludes a large portion of the m_ψ–β parameter space. The upper-left corner is also excluded due to the invisible decay of Z into WIMPs. The two horizontal lines correspond to the collider limits on four-fermion interactions with $\alpha' = 1$ and $\alpha' = \alpha$. (B) On the right, the parameter space of the singlet-Higgs mediated model. There are no collider constraints, and the $\sigma_{\rm el}$ constraint is much weaker.

Ref. [2], which is known to be essentially excluded by direct searches.

(b) Other particle physics constraints on this model are also highly dependent on m_V. For $m_V \gg m_Z$ and E, where E is the energy accessible in the collision, one can approximate V-exchange between SM particles as an effective current–current interaction,

$$\mathcal{L}_{\rm eff} = \frac{4\pi\alpha\kappa^2}{m_V^2 \cos\theta_W^2} J_\mu^Y J_\mu^Y. \tag{9}$$

We can then constrain the coefficient using limits on the corresponding effective four-lepton operator from LEP2 and Tevatron searches for "compositeness" [18,19]. In terms of the conventionally normalized coefficient, $4\pi/\Lambda_c^2$, no deviation is observed from the SM cross-section up to $\Lambda_c \sim 10$–15 TeV. Consequently, we arrive at the following collider constraint,

$$\Lambda_c < 10~\text{TeV} \quad \Longrightarrow \quad \beta < 0.3 \times \frac{\alpha'}{\alpha} \times \left(\frac{m_V}{500~\text{GeV}}\right)^2. \tag{10}$$

In order to plot this constraint in Fig. 3, we choose two representative values for the coupling constant e', one defining the perturbative regime, $\alpha' < 1$, and a more realistic line for $\alpha' = \alpha$.

(c) For $m_\psi < m_Z/2$, there is also an extra contribution to the invisible width of Z, namely $Z \to \psi\bar\psi$,

$$\Gamma_{Z\to\psi\bar\psi} = \frac{\alpha\beta}{3}\frac{m_Z^4}{(m_V^2 - m_Z^2)^2}\left(1 + \frac{2m_\psi^2}{m_Z^2}\right)\sqrt{m_Z^2 - 4m_\psi^2}. \tag{11}$$

Requiring this width to be less than 4 MeV [20] results in an additional constraint plotted in Fig. 3.

As is evident from the figure, the existing constraints already rule out WIMP masses up to 100 GeV with the exception of a small region in the vicinity of $m_\psi = m_V/2$ where there is resonant enhancement in the annihilation cross-section. Future LHC experiments, and the next generation of dark matter experiments, will provide much deeper probes into the parameter space of this model.

2.1.2. *The secluded WIMP regime*

The situation changes drastically if annihilation process (B) is kinematically allowed. The cross-section for WIMP annihilation into pairs of (unstable) V bosons is then given by

$$\sigma v = \frac{\pi(\alpha')^2}{m_\psi^2}\sqrt{1 - \frac{m_V^2}{m_\psi^2}}, \tag{12}$$

which together with (6) implies that in the limit $\beta \ll 1$ the correct dark matter abundance is achieved if

$$\alpha' \times \left(1 - \frac{m_V^2}{m_\psi^2}\right)^{1/4} \simeq 5 \times 10^{-3} \times \left(\frac{m_\psi}{500~\text{GeV}}\right). \tag{13}$$

This constraint is easily satisfied for a rather natural range for m_ψ and α'. Crucially, since large mixing is no longer required in this kinematic regime to ensure the correct annihilation cross-section, κ can be taken very small indeed. The only constraints one has to impose are that the decay of V (and also h') occur before the start of nucleosynthesis. By choosing $m_h' > m_V/2$, only the decays of V are sensitive to the mixing:

$$\Gamma_V \geqslant \text{s}^{-1} \quad \Longrightarrow \quad \kappa^2\left(\frac{m_V}{10~\text{GeV}}\right) \gtrsim 10^{-23}. \tag{14}$$

A tighter constraint would follow from requiring V decays to remain in thermal equilibrium, which would also ensure the initial thermal and chemical equilibrium for WIMPs, used in the derivation of the abundance formula (6),

$$\Gamma_V \geqslant \text{Hubble Rate}[T \simeq 0.05 m_\psi]$$
$$\implies \quad \kappa^2 \left(\frac{m_V}{10\ \text{GeV}}\right) \gtrsim 10^{-12} \left(\frac{m_\psi}{500\ \text{GeV}}\right)^2. \tag{15}$$

Although considerably tighter than (14), the constraint (15) does not change the main conclusion: in the limit $m_{\text{mediator}} < m_{\text{WIMP}}$, the WIMP sector can be secluded and neither collider nor underground searches impose any significant restrictions on the model. Moreover, the constraint (15) can be relaxed to (14) if some new UV physics ensures proper thermal contact between the dark matter and SM sectors at higher temperatures resulting in $T_{\text{DM}} \sim O(0.1) \times T_{\text{SM}}$ at the time of dark matter annihilation. For example, a dimension-8 operator of the form $V_{\mu\nu}^2 (F_{\mu\nu}^{\text{SM}})^2$ generated at a rather high energy scale would suit this purpose.

It is also worth noting that one could completely sever the remaining link to the SM, and set the mixing parameter κ to zero, *if* the $\text{U}(1)'$ gauge symmetry is not broken at all or broken only at very low energy scales. In this case, there will be an extra component to the energy density of the universe, namely the "dark radiation" associated with massless V-bosons. However, this energy density associated with V may be much smaller than that of the SM photons, because the decoupling of the dark sector may have occurred at a very early epoch, after which the photon temperature was effectively increased several-fold by input from the decay of numerous SM degrees of freedom. Although at first sight this model looks like a completely decoupled dark sector, the energy density in V could in principle be detected by highly sensitive next-generation CMB anisotropy probes at very small angular scales.

2.2. *Singlet scalar mediator*

An alternative class of models uses scalar mediators, discussed for example in several recent publications [10]. The simplest example consists of two additional fields, the singlet WIMP fermion ψ and a singlet mediator ϕ,

$$\begin{aligned}
&\mathcal{L}_{\text{WIMP+mediator}} \\
&\quad = \frac{1}{2}(\partial_\mu\phi)^2 - \frac{1}{2}m_\phi^2\phi^2 + \bar{\psi}(i\partial_\mu\gamma_\mu - m_\psi - \lambda_\psi\phi)\psi \\
&\qquad - \lambda_1 v\phi\left(H^\dagger H - \frac{v^2}{2}\right) - \lambda_2\phi^2\left(H^\dagger H - \frac{v^2}{2}\right) - V(\phi).
\end{aligned} \tag{16}$$

Here H is the SM Higgs doublet and v the corresponding vacuum expectation value, introduced in (16) for convenience. By a redefinition of the ϕ field, we can always set $\langle\phi\rangle = 0$. Should one pursue the minimalist approach, ϕ itself can be a dark matter candidate, which would then fix the value of λ_2. However, to demonstrate our point as simply as possible, we take the limit $\lambda_2 \to 0$. After electroweak symmetry breaking, the relevant dark matter Lagrangian then takes the following form,

$$\begin{aligned}
&\mathcal{L}_{\text{WIMP+mediator}} \\
&\quad = -\frac{1}{2}(\partial_\mu\phi)^2 + \bar{\psi}(i\partial_\mu\gamma_\mu - m_\psi - \lambda_\psi\phi)\psi - \lambda_1 v^2\phi h + \cdots,
\end{aligned} \tag{17}$$

where all interactions between SM fields and WIMPs are mediated by Higgs-singlet mixing.

The analysis of this model follows similar lines to the example above, so we will be rather brief. We first consider the case $m_\phi > m_\psi$ and choose $m_\phi = 400$ GeV to enable a clear comparison with the previous model. Working to lowest order in the mixing parameter, we have computed the annihilation cross-section taking into account $t\bar{t}$, ZZ, WW, hh and $b\bar{b}$ in the final state, following e.g. Ref. [9]. The cross-section has a leading p-wave contribution, which provides an additional suppression to $\langle\sigma v\rangle$.[1] Choosing the Higgs mass to be $m_h = 120$ GeV, and denoting the mixing parameter by β_h,

$$\beta_h \equiv \frac{\lambda_\psi^2 \lambda_1^2 v^4}{m_\phi^4}, \tag{18}$$

we can use the abundance constraint (6) to plot the dependence of β_h on the WIMP mass in Fig. 3(B). This plot also includes the experimental constraint on the nucleon-WIMP elastic scattering cross-section induced by h–ϕ scalar exchange. We take the SM value of the Higgs-nucleon coupling to be $g_{hNN} \simeq 300$ MeV. There are no significant collider constraints on this WIMP model, as all WIMP production has to occur via a real or virtual Higgs or Higgs-like bosons, and this is highly suppressed. Although quite constraining for low WIMP masses, in this scenario the direct search experiments cannot probe m_ψ above 50 GeV.

Again, the situation changes significantly if $m_\phi < m_\psi$, which allows for the WIMP sector to be secluded. The kinematically allowed pair-annihilation of WIMPs into two ϕ-scalars then simply imposes a constraint on a combination of λ_ψ, m_ϕ and m_ψ, but removes any constraints on λ_1. The decay of these scalars before BBN imposes a very relaxed requirement of $\lambda_1 \gtrsim 10^{-21}$. In practice, even taking $\beta_h \sim 10^{-5}$ would eliminate any chances for direct search experiments and/or colliders to probe the WIMP sector in such a model.

2.3. *Right-handed neutrino mediator*

In previous subsections the choice of metastable mediator, although relatively simple, was nonetheless exotic. Neither metastable vector or scalar particles are required by any known SM physics. There is, however, the distinct possibility of promoting right-handed neutrinos N_R, arguably the best motivated extension of the SM field content, to the role of metastable mediators with a dark matter sector. Indeed, if the right-handed neutrinos have Majorana masses of order the electroweak scale, their decay width into SM states such as left-handed neutrinos and the Higgs will be proportional to the square of the Yukawa coupling, i.e. in practice to the light neutrino masses. As is clear from the discussion of the previous subsections, $\Gamma_{N_R} \sim m_\nu \sim 0.1$ eV is in the right range for N_R to play the role of a metastable mediator. To complete this model, we choose the secluded sector to have one additional fermion N' and a bo-

[1] The presence of a pseudoscalar coupling, $\phi\bar{\psi}i\gamma_5\psi$, would open up the s-channel.

son S:

$$\begin{aligned}
&\mathcal{L}_{\text{WIMP+mediator}} \\
&\quad = \frac{1}{2}(\partial_\mu S)^2 - \frac{m_S^2}{2}S^2 + \bar{N}' i\partial\!\!\!/ N' - \frac{m_{N'}}{2}N'^T N' \\
&\qquad + \bar{N}_R i\partial\!\!\!/ N_R - \frac{m_{N_R}}{2}N_R^T N_R - \frac{\lambda}{2}S^2 H^\dagger H \\
&\qquad - \left[Y_\nu \bar{L} H N_R - Y_{N'} S N_R^T N' + (\text{h.c.})\right].
\end{aligned} \tag{19}$$

This Lagrangian possesses a Z_2 symmetry which allows only even numbers of N' or S at each vertex. Alternatively, one may complexify S and N' and introduce a new global charge that would ensure the same property.

A viable secluded dark matter model results from the choice $m_S, m_{N'} > m_{N_R}$. If $m_S > m_{N'}$, then S-mediated annihilation $N' + N' \to N_R + N_R$ will ensure the right WIMP abundance of N' given an appropriate choice of $Y_{N'}$. Notice that a scalar coupling λ as small as 10^{-8} is sufficient to keep the dark sector in thermal/chemical equilibrium prior to freeze-out [9]. For $m_S < m_{N'}$ the roles are reversed, and N'-mediated annihilation allows for S to be dark matter. In either case, the dark matter candidate, either N' or S, is secluded from direct observational probes.

2.4. Strong non-Abelian interactions in the secluded sector

Models with a non-Abelian gauge group G' in a hidden sector can also be secluded and may be of interest in the context of dark matter models with strong self-interactions [15] if the scale is relatively light. The gauge bosons of G' cannot couple directly to the SM because of gauge invariance. Therefore, in this case the mediators must be charged under both gauge groups. The simplest example of this kind is given by

$$\begin{aligned}
&\mathcal{L}_{\text{WIMP+mediator}} \\
&\quad = -\frac{1}{4}\left(G^a_{\mu\nu}\right)^2 + \sum_f \bar{f}\left(i\gamma_\mu D_\mu^{\text{SM,hid}} - m_f\right) f \\
&\qquad + \sum_\psi \bar{\psi}\left(i\gamma_\mu D_\mu^{\text{hid}} - m_\psi\right)\psi,
\end{aligned} \tag{20}$$

where $G^a_{\mu\nu}$ is the non-Abelian field strength in the hidden sector, while ψ and f are fermions, with ψ charged only under G', while the fermions f play the role of mediators and are charged under both G' and the SM gauge group. This field content will necessarily have to satisfy anomaly cancelation constraints. We further assume that the confinement scale of G' is comparable to or larger than the electroweak scale, $\Lambda' > v$.

The Lagrangian (20) allows the construction of a secluded WIMP model, in which both mediators and dark matter are composite. If all masses in the f-sector are large, these fields can be integrated out leading to non-renormalizable interactions between the two sectors of the form $\frac{1}{m_f^4}(G^a_{\mu\nu})^2(F^{\text{SM}}_{\mu\nu})^2$. The phenomenology of such terms, in connection with metastable dark matter composed of hidden-sector glueballs was considered in Ref. [21]. The ψ-containing baryons B' and antibaryons are viable dark matter candidates if the nonrenormalizable terms leading to their decays are forbidden. The mass of an exotic "meson" ($M' \sim \bar{\psi}\psi$) would be smaller than the baryon mass, and the annihilation process $B'\bar{B}' \to M'\bar{M}'$ therefore open. Since this cross-section is not expected to have additional suppression other than that provided by the mass scale of the exotic baryons, the abundance constraint would require that the dark matter mass be above 10 TeV. The mesons M' as well as the glueballs of the hidden group would be unstable with respect to decay into the SM fields, with total widths controlled by the combination Λ'^9/m_f^8 [21]. For large m_f, i.e. $m_f \gg \Lambda'$, this width can be exceptionally small, e.g. of order the Hubble scale during $B\bar{B}$ freeze-out. Explicit models with exotic baryons as dark matter have previously been constructed within the framework of gauge-mediated supersymmetry breaking [22,23].

The opposite limit, $m_f \ll \Lambda'$, would provide a regime where the WIMP sector could in principle be probed because the mesons built from f would then be of electroweak scale, and will interact both with the SM and with WIMPs in the form of exotic baryons. Although the direct detection signal might again be rather low, the exotic f-containing mesons could conceivably be produced in the next generation of colliders [25].

3. MeV-scale dark matter and mediators

The choice of a relatively light mediator is perhaps the only viable possibility for WIMP masses to lie well below the Lee–Weinberg window, and close to the MeV-scale [24]. This mass range has some interesting consequences, including the speculative possibility of linking the unexpectedly strong and uniform emission of 511 keV photons from galactic center [7] to positrons created by the annihilation $O(\text{MeV})$-scale dark matter [6]. To date, much of the model-building in this direction has concentrated on utilizing an additional U(1) gauge group, under which both the Standard Model and the dark sector are (disproportionately) charged: a small charge for the Standard Model fermions, and a larger charge for dark matter [24,26].

The most natural anomaly-free quantum number to gauge is $B - L$ [27], in which case the coupling of the additional U(1) gauge bosons to charged fermions is necessarily vector-like. The absence of an axial vector current allows this scenario to escape strong constraints from atomic parity violation, and from flavor-changing decays of K and B mesons [26]. However, there is a price to pay as the coupling to neutrinos creates a problem with the energetics of supernovae. During the explosion, the MeV-scale WIMP is thermalized and coupled to neutrinos too strongly, suppressing the energy of the emitted neutrinos and making the observed SN1987A signal highly unlikely [28]. There are two ways to escape this problem: taking the WIMP mass in excess of 10 MeV, or forbidding couplings to neutrinos [28]. Unfortunately, the first option does not work, because the shape of the 511 keV line [29], along with the γ-ray spectrum in the MeV region [30], do not allow the mass of annihilating particles to be in excess of 3–5 MeV. The second option, i.e. no coupling to neutrinos, requires abandoning the initial assumption of gauging $B - L$.

In this context, it is easy to see that the secluded models of Sections 2.1 and 2.2 with vector and scalar mediators can solve certain model building problems for MeV-scale dark matter. Indeed, neither of these mediation mechanisms lead to any significant coupling to neutrinos. In particular, for the U(1)$'$ mediator the ψ–ν scattering amplitude is necessarily suppressed by the Z-boson mass, and since the mixing parameter κ is smaller than 1, the ψ–ν scattering cross-section will be below the typical weak-scale value. This motivates closer inspection of these two secluded models in the MeV mass range, in order to determine if they do indeed represent viable MeV-scale WIMP scenarios. In what follows we will make use of two relations that generalize results of [26,31], namely for the abundance.

$$\frac{\Omega_X}{\Omega_{\rm DM}} \simeq (2\text{–}4)\frac{10^{-36}\ {\rm cm}^2}{\langle\sigma v\rangle_c}, \tag{21}$$

and the 511 keV flux,

$$\frac{\Phi_{511,X}}{\Phi_{511,\rm total}} \sim \frac{N_{e^+}\langle\sigma v\rangle_g}{10^{-40}\ {\rm cm}^2} \times \left(\frac{1\ {\rm MeV}}{m_X}\right)^2 \times \left(\frac{\Omega_X}{\Omega_{\rm DM}}\right)^2, \tag{22}$$

where m_X is the mass of the dark matter candidate X (with $0.5\ {\rm MeV} < m_X \lesssim 3\text{–}5\ {\rm MeV}$), $\Omega_X/\Omega_{\rm DM} \leqslant 1$ is the contribution of X to the total dark matter energy density, and the ratio $\Phi_{511,X}/\Phi_{511,\rm total}$ is the contribution of X-annihilation to the total 511 keV γ-ray flux from the galactic center. In (22), N_{e^+} is the positron multiplicity per annihilation event. Since the neutrino couplings are negligibly small, and direct annihilation to photons is also suppressed [32], in these models e^+e^- *is* the dominant annihilation mode, and $N_{e^+} = 1$ for $m_{\rm WIMP} < m_{\rm mediator}$, and $N_{e^+} = 2$ otherwise. The subscript on the annihilation cross-section, $\langle\sigma v\rangle_c$ or $\langle\sigma v\rangle_g$, denotes the type of averaging: $\langle\cdots\rangle_c$ implies thermal averaging for cosmological freeze-out; while $\langle\cdots\rangle_g$ refers to the average over the galactic velocity distribution for the X-particles. We should caution the reader that the relation (22) is only valid to within one to two orders of magnitude, due to the uncertainty in the dark matter number density in the galactic core.

(a) U(1)$'$-*mediator*, $m_X > m_V$: Irrespective of whether the dark matter is a fermion or a scalar, annihilation to two V bosons proceeds in the s-channel, and $\langle\sigma v\rangle_c \simeq \langle\sigma v\rangle_g$. This immediately rules out MeV-scale X particles as the dominant component of dark matter, since $\Omega_X/\Omega_{\rm DM} \simeq 1$ will ensure that the 511 keV γ-ray flux is significantly overproduced, $\Phi_{511,X}/\Phi_{511,\rm total} \sim 10^4$. Alternatively, one can tune the coupling α' to be much smaller than the SM gauge couplings, $\alpha' \sim 2\times 10^{-3}\times\alpha$, and satisfy the 511 keV flux constraint with $\Omega_X/\Omega_{\rm DM} \sim 2\times 10^{-5}$. This scenario is reminiscent of the decaying sterile neutrino [33], where a very small contribution of sterile neutrinos to the total dark matter budget can provide the requisite flux. The dominant component of dark matter should of course come from other sources, which renders this model incomplete.

(b) U(1)$'$-*mediator*, $m_X < m_V$: In this regime, the choice of X as a scalar charged under $U(1)'$ is preferred [24], because of the p-wave annihilation that leads to $\langle\sigma v\rangle_g \sim 10^{-5}\times\langle\sigma v\rangle_c$. It then appears entirely possible to satisfy both (21) and (22) with $m_X \sim O({\rm MeV})$, $\Omega_X/\Omega_{\rm DM} \simeq 1$, and a mixing parameter,

$$\beta \sim {\rm few}\times 10^{-6}\times\left(\frac{m_V}{10\ {\rm MeV}}\right)^4. \tag{23}$$

This value for the mixing parameter would be natural for example if induced radiatively at a high scale by a state charged both under the SM U(1) and the extra U(1)$'$. As mentioned before, this model does not pose any problems with respect to the supernova signal, and does not presuppose any hierarchy of gauge couplings as α' can be taken of order α. Therefore, this model appears the most natural candidate for MeV-scale secluded dark matter, having the chance to explain the 511 keV line from the galactic center.

(c) ϕ-*mediator*, $m_X > m_\phi$: In this scenario, it is advantageous to have a fermionic dark matter candidate ψ with scalar (rather than pseudoscalar) couplings to ϕ. The annihilation $\psi\psi \to \phi\phi$ proceeds in the p-wave and can always be tuned to the required level with a typical choice $\lambda_\psi \sim 10^{-6}$. Since $m_\psi \sim$ few MeV, this value of the Yukawa coupling is natural. The subsequent decay of ϕ due to mixing with the Higgs is highly suppressed by the electron Yukawa coupling,

$$\Gamma_\phi \sim \left(\frac{\lambda_1 v^2}{m_h^2}\right)^2 \times \left(\frac{m_e}{v_{\rm EW}}\right)^2 \times \frac{m_\phi}{8\pi} \gtrsim {\rm s}^{-1}$$
$$\Longrightarrow \quad \left(\frac{\lambda_1 v^2}{m_h^2}\right)^2 \gtrsim 10^{-8}. \tag{24}$$

The naturalness requirement for the ϕ-mass would impose a significant constraint here. If we consider the contribution from Higgs mixing in (17), $\lambda_1 v/m_h \lesssim m_\phi/v$, this clearly favors a long ϕ-lifetime (~ 1 s) and a small mixing parameter. Even then, one must ensure that the "missing energy" decay $K^+ \to \pi^+ + \phi$ is within the allowed range. At the quark level, the amplitude for the process is given by a Higgs penguin (see, e.g. [34]):

$$\mathcal{L}_{\rm eff} = \left(\frac{\lambda_1 v^2}{m_h^2}\right)\frac{3g_W^2 m_s m_t^2 V_{td}V_{ts}^*}{64\pi^2 m_W^2 v}\bar{d}_L s_R \phi + ({\rm h.c.}), \tag{25}$$

leading to the (non-SM) missing energy decay,

$$\Gamma_{K\to\pi+\phi-{\rm mediator}} \simeq \left(\frac{\lambda_1 v^2}{m_h^2}\right)^2\left(\frac{3m_t^2 V_{td}V_{ts}^*}{16\pi^2 v^2}\right)^2\frac{m_K^3}{64\pi v^2}. \tag{26}$$

Requiring that this width not exceed the observed missing energy decay branching ratio ${\rm Br} = 1.5^{+1.3}_{-0.9}\times 10^{-10}$ [35] associated with the SM process $K^+ \to \pi\nu\bar{\nu}$, results in the following constraint on ϕ–h mixing:

$$\left(\frac{\lambda_1 v^2}{m_h^2}\right)^2 < 2\times 10^{-7}. \tag{27}$$

This cuts out a significant part of the parameter space, but together with (24) still leaves a relatively narrow interval for the mixing parameter, 10^{-7}–10^{-8}, where the model survives all constraints (although not without a modest amount of fine-tuning of the mediator mass) and thus can be the dominant dark matter component while still accommodating the positron signal through a combination of annihilation and decay.

The constraints remain essentially the same for a pseudoscalar coupling of ϕ to the fermion ψ, if the Higgs sector in SM is assumed to be minimal, in which case the mixing constant λ_1 is CP-violating. The additional processes: s-wave annihilation $\psi\psi \to e^+e^-$ through a virtual ϕ, and also $\psi\psi \to \phi\phi\phi$ if kinematically allowed, are too weak in comparison with the p-wave annihilation $\psi\psi \to \phi\phi$ to affect the constraints discussed above. In principle, with an extended Higgs sector, ϕ could also mix in a CP-conserving way with the physical CP-odd Higgs scalar(s) and in addition could have an enhanced coupling to electrons, relative to the top quark, thus significantly relaxing the constraints on the parameter space in comparison with Eqs. (24) and (27).

(d) *ϕ-mediator*, $m_X < m_\phi$: In this case the annihilation cross-section is suppressed by $(m_e/v)^2 \sim 4 \times 10^{-12}$ and the relation (21) would require a mixing parameter, β_h, of order one. Such a choice would involve a gross violation of naturalness, and would also lead to an unacceptably large missing energy signal in K and B decays [34]. We therefore conclude that this option is not viable.

4. Concluding remarks

Secluded WIMP dark matter appears to be a generic possibility, as rather minimal model-building choices lead to viable WIMPs interacting with metastable mediators. These mediators could be either elementary or composite, and we have constructed explicit models with scalar, vector and fermion mediation to the Higgs, hypercharge gauge boson, and light neutrino sectors respectively of the SM. In the latter case there is a natural choice for the mediator, namely a right-handed electroweak-scale neutrino.

Despite existing as a thermal relic, with weak-scale annihilation, none of the secluded WIMP models constructed here lead to appreciable signals in underground detectors or register as a missing energy channel in collider experiments. Nonetheless, these models are subject to indirect constraints related to γ-rays caused by WIMP annihilation, e.g. in the galactic center. Therefore, one of the main conclusions of this Letter is the complete complementarity of direct and indirect efforts for detecting WIMP dark matter.

Two of the WIMP models constructed here allow a rescaling down to the MeV mass range, motivated by the intriguing connection to the galactic 511 keV line. The secluded models considered here do not couple dark matter to neutrinos, and therefore avoid problems with the suppression of SN1987A signal. We have shown that models with an additional $U(1)'$ and kinetic mixing to the hypercharge gauge boson circumvent some problems of scenarios where SM fields carry an additional gauge quantum number, and appear to be free from unnaturally small parameters. The main limitations come instead from the dark matter energy density and from indirect constraints related to the galactic γ-ray spectrum in the MeV region.

Note added

As this Letter was being finalized, we became aware of a recent preprint [36] that also deals with $U(1)'$ models of MeV-scale dark matter with kinetic mixing, and thus has some overlap with the discussion in Sections 3(a) and (b).

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