

Solar gamma rays powered by secluded dark matterBrian Batell,¹ Maxim Pospelov,^{1,2} Adam Ritz,² and Yanwen Shang¹¹*Perimeter Institute for Theoretical Physics, Waterloo, Ontario, N2J 2W9, Canada*²*Department of Physics and Astronomy, University of Victoria, Victoria, British Columbia, V8P 1A1 Canada*

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Secluded dark matter models, in which weakly interacting massive particles annihilate first into metastable mediators, can present novel indirect detection signatures in the form of gamma rays and fluxes of charged particles arriving from directions correlated with the centers of large astrophysical bodies within the Solar System, such as the Sun and larger planets. This naturally occurs if the mean free path of the mediator is in excess of the solar (or planetary) radius. We show that existing constraints from water Cerenkov detectors already provide a novel probe of the parameter space of these models, complementary to other sources, with significant scope for future improvement from high angular resolution gamma-ray telescopes such as Fermi-LAT. Fluxes of charged particles produced in mediator decays are also capable of contributing a significant solar system component to the spectrum of energetic electrons and positrons, a possibility which can be tested with the directional and timing information of PAMELA and Fermi.

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I. INTRODUCTION

The search for weakly interacting massive particles (WIMPs) as a component of nonbaryonic dark matter has become a focal point of modern particle physics [1]. There are several complementary experimental and observational approaches to WIMP detection [2]. Direct detection experiments probe the terrestrial scattering of WIMPs with nuclei and typically require low radiation environments to keep backgrounds under control. High-energy colliders such as the Tevatron and the LHC offer the possibility of producing WIMPs and measuring their properties in the laboratory, provided the challenging missing energy signatures can be disentangled. Indirect searches for dark matter annihilating into gamma and cosmic rays in the galactic halo are also promising, although susceptible to various, often uncertain, astrophysical backgrounds. Finally, neutrino telescopes such as Super-Kamiokande and Ice Cube can search for indirect evidence of the annihilation of WIMPs captured in the core of the Sun and the Earth, in the form of an observable muon signature arising from neutrino charged current scattering in the detector. While the latter two examples are well-known indirect signatures for any thermal relic WIMP dark matter candidate, more generic WIMPs forming part of a larger dark sector can lead to further novel signatures. In this paper, we demonstrate that models of secluded dark matter [3] present an additional observational possibility: high-energy gamma rays and charged particles arriving from a direction tightly correlated with the centers of the Sun, Earth and other planets. Such novel signatures can be effectively probed with the powerful new generation of gamma-ray telescopes.

The primary feature of secluded models of dark matter [3] is a two-stage dark matter annihilation process: WIMPs

annihilate first into metastable mediators, which subsequently decay into standard model (SM) states. This breaks the more-or-less rigid link between the size of the WIMP annihilation and WIMP-nucleus scattering cross sections. It has been shown that a small mass for the mediator allows for new phenomenological possibilities in the form of enhanced WIMP annihilation at small velocities [4,5] that may help to explain various astrophysical anomalies, e.g. the positron excess observed by PAMELA above 10 GeV [6] and perhaps the unexpectedly hard electron spectrum observed by Fermi above a few hundred GeV [7]. Furthermore, a relatively small mediator mass kinematically removes heavy SM particles from the final state [4,5], reconciling these effects with the absence of any enhancement in the cosmic ray antiproton signal [8]. The lifetime of the mediator is essentially a free parameter, limited only by the big bang nucleosynthesis bounds of $\tau \lesssim 1$ s. If this lifetime is rather long, the decay of the mediator will occur a long distance away from the point of the original WIMP annihilation. Denoting the WIMP particle χ and the mediator particle V , assuming $\chi\chi \rightarrow 2V$ as the main annihilation channel, and taking $m_V \ll m_\chi$, we arrive at the following estimate for the mediator travel distance:

$$L = c\tau_V\gamma_V = 3 \times 10^6 \text{ km} \times \frac{\tau_V}{0.01 \text{ s}} \times \frac{\gamma_V}{10^3}. \quad (1)$$

Such large boosts $\gamma_V = m_\chi/m_V$ are easily achieved if the dark matter mass is near the electroweak scale and the mediator mass is below a GeV. With regard to the annihilation of WIMPs captured within the Sun, one can see that this distance may very well exceed the solar radius ($R_\odot = 6.96 \times 10^5$ km) in which case, unlike conventional WIMPs, most of the decay products will not be absorbed. For somewhat shorter lifetimes, the interesting possibility emerges of mediators produced by WIMP annihilation in

the center of the Earth decaying directly into charged particles within neutrino telescopes. A schematic illustration of this new indirect mechanism of probing secluded WIMP models is shown in Fig. 1.

A sub-GeV mass mediator will decay predominantly into light states, such as pions, muons, electrons, gammas and neutrinos. As they are produced in the decays of highly boosted mediators, these decay products will be tightly correlated with the direction to the original WIMP annihilation point, with a typical angular size $\theta_V \sim 1/\gamma_V$. While for charged particles this correlation can be reduced by the magnetic fields encountered on the way to the detector, the directionality of gammas and neutrinos is unaltered. Modern gamma-ray telescopes enjoy an angular resolution much better than a degree, which may be exploited to enhance the gamma-ray signal-to-background. In particular, a notable source of background to this signature is the generation of gamma rays via cosmic rays impinging on the Sun (see, e.g. [9,10]). Such background gammas will typically display a much softer spectrum than gamma rays from mediator decays and will have no specific correlation with the solar center where most of the secluded WIMP annihilation is expected to take place. The main gain in sensitivity in detecting gammas from the decays of metastable mediators, as compared to a more conventional search for a highly-energetic neutrino signal, may come from the increase in efficiency. While the detection of multi-GeV neutrinos requires their conversion to muons, which means a loss in efficiency of around 10 orders of magnitude, the efficiency of detecting gammas created outside of the solar radius can be order one.

In addition to gamma rays, many secluded WIMP models with relatively light mediators are destined to produce a significant fraction of leptons in the final state, and thus we are naturally led to the question of whether mediator decays outside the solar/planetary radii are capable of contributing significantly to the fluxes of electrons and positrons seen by PAMELA and Fermi. It is tempting to

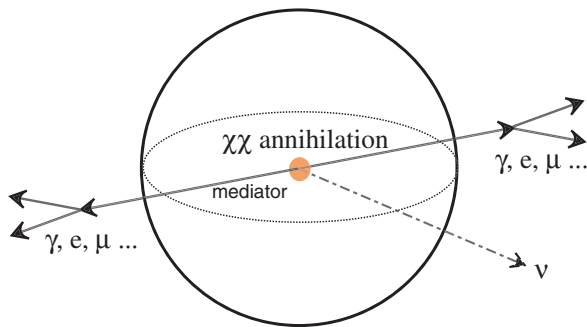


FIG. 1 (color online). A schematic illustration of the new indirect detection signature of secluded WIMPs captured in the solar core, annihilating to metastable mediators and leading to an electromagnetic flux: γ , e^\pm , μ^\pm , \dots . Sensitivity to conventional WIMPs arises only through annihilation to neutrinos.

pursue the notion that these anomalies may in fact have a local origin in the Solar System, powered by the annihilation of secluded dark matter trapped within solar bodies. This is an intriguing possibility, as it would offer a new dark matter interpretation of these signatures that does not rely on *galactic* WIMP annihilation, which generally requires a significant boost factor in the annihilation cross section. In this context, a less extreme hierarchy between the WIMP and mediator mass may be feasible, and this idea is akin to using local substructure to enhance the charged particle flux, albeit with a source (the Sun) which is extremely local and well-understood. These electrons and positrons, like the gamma rays discussed above, should be strongly correlated with the center of the Sun, particularly in the high-energy range where the effect of the magnetic fields of the Sun and Earth is less significant. Such a hypothesis appears straightforwardly falsifiable using directional and timing data from PAMELA and Fermi. Finally, if charged particles are the primary signature, this immediately leads to a minimal gamma-ray flux from the associated 3-body decays that is also directionally correlated with the center of the Sun, and thus testable by Fermi-LAT.

In this paper we analyze the feasibility of detecting electromagnetic particles, γ , e^\pm , \dots , arising from the delayed decays of metastable mediators produced through WIMP annihilation in the deep interior of the Sun and other planetary bodies within the Solar System. Our primary focus is on the capability of modern gamma-ray telescopes to search for these local annihilation signatures and the possibility of distinguishing the signal gammas from the solar backgrounds produced by cosmic rays. We demonstrate that in certain corners of the secluded dark matter parameter space, the gamma-ray signature from the center of the Sun is potentially the most sensitive probe, superior to direct detection and other indirect signatures. This is especially the case for models in which WIMP-nucleus scattering is dominated either by spin-dependent interactions or proceeds through an inelastic transition to a nearby excited state. We also analyze the prospects for detecting gamma rays generated by dark matter annihilations inside Jupiter using atmospheric Čerenkov gamma-ray detectors, as well as the potential for observing pairs of upward going muons in neutrino telescopes generated by the decay of mediators to muons. Finally we offer some preliminary speculations concerning an interpretation of PAMELA's rising positron fraction originating within the Solar System, and discuss the obstacles to this identification as well as "smoking gun" directional and temporal signatures which can confirm or rule out such an interpretation.

The rest of this paper is organized as follows. In the next section we present the WIMP trapping rates and general formulas for calculating the WIMP-powered gamma-ray flux, including angular and energy distributions. Section III

contains estimates of the gamma-ray signature in two variants of secluded models, with pseudoscalar (“axion”)- and vector-mediation. We conclude in Sec. IV with a discussion of future prospects for improving sensitivity to these dark matter models via observation of gamma rays and charged particles in the Solar System.

II. CAPTURE AND DELAYED ELECTROMAGNETIC DECAYS OF MEDIATORS

In this section we derive the general formulas for the expected gamma-ray fluxes generated by metastable mediator decays. We will also consider the case of charged particle fluxes and comment on the expected effect of the magnetic fields of the Sun and Earth.

A. Solar capture and γ -flux

We begin by providing the trapping efficiency of WIMPs inside the Sun, assuming elastic scattering of WIMPs on nuclei [1,11,12] and normalizing unknown quantities on some fiducial values:

$$C_{\odot} \simeq 1.3 \times 10^{21} \text{ s}^{-1} \times \left(\frac{100 \text{ GeV}}{m_{\chi}} \right) \sum_N f_N \left(\frac{\sigma_N^{\text{SD}} + \sigma_N^{\text{SI}}}{10^{-42} \text{ cm}^2} \right) S(m_{\chi}/m_N) F_N(m_{\chi}), \quad (2)$$

where the sum runs over the nuclei N present in the Sun, f_N denotes the fractional abundance relative to hydrogen, σ_N^{SD} (σ_N^{SI}) is the spin-dependent (spin-independent) scattering cross section and we have used the standard values $\rho_{\chi} = 0.3 \text{ GeV cm}^{-3}$ and $\bar{v} = 270 \text{ km/s}$ for the local WIMP density and velocity dispersion. The function $S(x) = [A(x)^{3/2}/(1 + A(x)^{3/2})]^{2/3}$, where $A(x) = (3/2) \times [x/(x-1)^2](v_{\text{esc}}/\bar{v})^2$, with $v_{\text{esc}} \simeq 1156 \text{ km/s}$ an “effective” escape velocity, is a kinematic suppression factor, while $F_N(m_{\chi})$ determines additional suppression from a nuclear form-factor. This approximate formula holds for both spin-independent and spin-dependent scattering, but it is important to bear in mind that coherent scattering in the former case means that heavier nuclei, such as He, O, and Fe tend to dominate the capture rate despite their reduced abundance. The lack of nuclear coherence in spin-dependent scattering means that accounting for scattering off hydrogen is generally sufficient. More general formulas may be found in Ref. [11].

For most models the trapping rate in the Sun determines the overall annihilation rate, as the two processes, trapping and annihilation, are usually in dynamical equilibrium. With an S -wave annihilation rate of order 1 pb as dictated by the relic abundance, a per-nucleon scattering cross section of order 10^{-48} cm^2 or larger is generally sufficient for thermalization to occur over the lifetime of the Sun. Then the trapping rate C_{\odot} and the probability of a single annihilation event to produce gamma quanta reaching the

Earth, which we denote as P_{γ} , determines the overall flux of solar WIMP-generated gamma rays at the Earth’s location:

$$\begin{aligned} \Phi_{\gamma\odot} &= \frac{1}{2} \times \frac{C_{\odot} P_{\gamma}}{4\pi (\text{A.U.})^2} \\ &= 1.8 \times 10^{-7} \text{ cm}^{-2} \text{ s}^{-1} \times \frac{C_{\odot} P_{\gamma}}{10^{21} \text{ s}^{-1}}. \end{aligned} \quad (3)$$

This equation is correct if the particles that produce the final photons are highly boosted along the radial direction of the Sun, and should be modified by another factor of $1/2$ if those particles travel very slowly and emit photons isotropically as the Sun is effectively opaque and photons emitted on the back side of the Sun are reabsorbed. We are only interested in the former case in the current discussion. The master formula (3) suggests that observable gamma-ray fluxes are indeed possible, provided that P_{γ} is not too small. In conventional neutralino-like WIMP scenarios, this probability is in fact negligible, as the only practical way of producing gamma rays is via high-energy neutrinos interacting with the outer layer of the solar material,¹ and in this case the probability P_{γ} is very small. In contrast, secluded WIMPs offer the possibility of maximizing P_{γ} , which arises as the product of: (i) the probability P_{out} that the mediator particle decays outside the solar radius; (ii) the probability Br_{γ} of producing a gamma quantum in the decay of the mediator; and (iii) the probability Br_V of producing the mediator particle in the annihilation process:

$$P_{\gamma} = g \times P_{\text{out}} \times \text{Br}_{\gamma} \times \text{Br}_V. \quad (4)$$

Here, $g = g_{\gamma} g_V$ where g_{γ} and g_V are the multiplicities of the photon and the mediator particle V being produced along the chain of reactions. If the loss of V due to rescattering inside the Sun can be neglected, the mediator escape probability P_{out} is well approximated by

$$P_{\text{out}} \simeq \exp\{-R_{\odot}/(c\tau\gamma)\}, \quad (5)$$

while the branching ratios Br_V and Br_{γ} and the multiplicities are model-dependent. For example, Br_{γ} may be close to 1 in some secluded models, or more commonly may lie in the 10^{-3} – 10^{-2} range for models where gammas are produced as radiation accompanying the decay to charged particles.

The general relations (2) and (3) use elastic cross sections normalized to 10^{-42} cm^2 , in conflict with the current bounds on spin-independent elastic WIMP-nucleon scattering [15], but well below the bounds on spin-dependent scattering [16]. Therefore, the best sensitivity to secluded WIMPs will occur for models where nucleon scattering is predominantly spin-dependent. Alternatively, scattering

¹Earlier claims of enhanced neutralino annihilation immediately outside of the solar radius [13] were not confirmed by subsequent studies [14].

may be predominantly inelastic $\chi_1 \rightarrow \chi_2$, with a small mass gap between the two WIMP states relaxing the most stringent constraints on spin-independent σ_p . We will consider examples in both these categories in Sec. III.

B. Angular and spectral distributions

We now turn to the angular and spectral distributions of gamma quanta. We will comment at the end of this section on how these distributions may be generalized for charged particles, such as e^\pm , for which the propagation effects of the magnetic fields of the Sun and Earth must be taken into account. For simplicity we shall assume a decay of the mediator into a pair of gammas, which are monochromatic and isotropic in the rest frame of the mediator. The Lorentz boost of the mediator then determines both the angular and energy resolution. In the absence of the boost, $\gamma \sim 1$, one would expect the Sun to acquire a constant surface brightness in gamma rays, assuming that the decays happen not far from the solar radius, although its side would appear brighter due to simple geometrical reasons. However, in the most interesting case of large boosts, $\gamma > 10^2$, the gamma quanta in the decay products are emitted within an angle $\theta \sim 1/\gamma$ from the original direction of the mediator. For γ as large as 10^3 , the majority of gamma rays would come from an angular spot in the sky smaller than the solar radius, and such a tight correlation is observable with modern gamma-ray telescopes. Also, the gamma energy spectrum would be peaked in the direction of the solar center, although this effect is significantly smoothed out by the finite angular resolution of the detectors. In this subsection, we illustrate these angular distributions and energy spectra for the case of $m_\chi = 1$ TeV WIMPs annihilating into two metastable mediators, that further decay into 2γ , so that the multiplicity factors in Eq. (4) are $g_V = g_\gamma = 2$.

In the rest frame of the intermediate particle V , photons are emitted isotropically with a momentum distribution $f'(p') = \frac{2}{\pi m_V^2} \delta(p' - m_V/2)$, where $p = |\vec{p}|$. After a Lorentz transformation, the distribution function $f(\vec{p})$ measured in the observer's frame becomes

$$f(p, \alpha) = \gamma(1 - \beta \cos\alpha) f'[\gamma(1 - \beta \cos\alpha)p], \quad (6)$$

where $\beta \equiv v/c$, and α is the angle between the momentum of the outgoing photon and that of the particle V measured in the observer's frame. Consequently, photons are emitted with an angular distribution given by $1/[4\pi\gamma^2(1 - \beta \cos\alpha)^2]$. The density of the photons arriving at the detector with an incoming angle θ whose momenta are within a small solid angle $d\Omega$ is calculated by integrating the particle density of V along the line of sight weighed by an appropriate angular distribution factor:

$$\frac{d\Phi(\theta)}{d\Omega} = \frac{\cos\theta \text{Br}_\gamma}{2\pi\tau} \int_D d\cos\alpha \frac{n(l \sin\theta \csc\alpha) l \sin\theta}{\gamma^3 \sin^3\alpha (1 - \beta \cos\alpha)^2}. \quad (7)$$

Here $l \approx 1$ A.U. is the Sun-detector distance and $n(r)$ is the number density of V at distance r from the center of the Sun. Given our assumptions,

$$n(r) = \frac{C_\odot \cdot \text{Br}_V}{4\pi v r^2} \exp\{-r/(v\tau\gamma)\}, \quad (8)$$

where $v \simeq c(1 - 1/(2\gamma^2))$ is the velocity of V .

A few words about the range of the integral for $d\cos\alpha$, which we have denoted as D above, are in order. If the Sun were transparent, D would simply be the domain $[-1, +1]$. In the current case, however, the Sun is effectively opaque since photons traveling through the interior are instantaneously absorbed or degraded in energy through their interactions with solar material. Thus one should simply discard those photons from consideration, which leads to a lower bound for the integral at $\cos\alpha = \sqrt{1 - (l \sin\theta)^2/R_\odot^2}$ whenever $l \sin\theta < R_\odot$. Consequently,

$$D = \begin{cases} [\sqrt{1 - (l \sin\theta)^2/R_\odot^2}, +1], & l \sin\theta < R_\odot; \\ [-1, +1], & l \sin\theta \geq R_\odot. \end{cases} \quad (9)$$

The lower integration limits are not important for $\gamma \gg 1$, as almost all the photons detected come from within a small angle $\sim 1/\gamma$ and the flux is completely negligible when $\theta \approx \theta_\odot$. We illustrate $d\Phi/d\Omega$ for a representative value of $\gamma = 10^3$ in Fig. 2. Given a realistic angular resolution, it is clear that all events associated with the decays of highly boosted mediators will be concentrated in the angular bin covering the solar center.

The energy distribution of Φ follows from a similar integral. In general, we can write the full differential distribution in the form,

$$\begin{aligned} \frac{d^2\Phi(\theta, p)}{p^2 dp d\Omega} &= \frac{l \sin\theta \cos\theta \text{Br}_\gamma}{\gamma\tau} \\ &\times \int_D d\cos\alpha \cdot \frac{f(p, \alpha) n(l \sin\theta \csc\alpha)}{\sin^3\alpha} \\ &= \frac{l \sin\theta \cos\theta \text{Br}_\gamma}{\tau} \int_D \frac{d\cos\alpha}{\sin^3\alpha} \cdot (1 - \beta \cos\alpha) \\ &\times f'[\gamma p(1 - \beta \cos\alpha)] n(l \sin\theta \csc\alpha), \end{aligned} \quad (10)$$

where D is the same domain explained above. Specializing to the case of two-body decays of V to photons, we have

$$\begin{aligned} \frac{d^2\Phi(\theta, p)}{p^2 dp d\Omega} &= \frac{(\gamma^2 - 1) \sin\theta \cos\theta l \text{Br}_\gamma}{2\pi\gamma\tau(2\gamma p_0 p - p_0^2 - p^2)^{3/2}} \cdot \frac{p}{p_0} \\ &\cdot n\left(\frac{\sqrt{\gamma^2 - 1} \sin\theta l p}{\sqrt{2\gamma p_0 p - p_0^2 - p^2}}\right), \end{aligned} \quad (11)$$

where $p_0 \equiv m_V/2 = m_\chi/(2\gamma)$ is the photon energy measured in the rest frame of V . When $l \sin\theta \geq R_\odot$, this

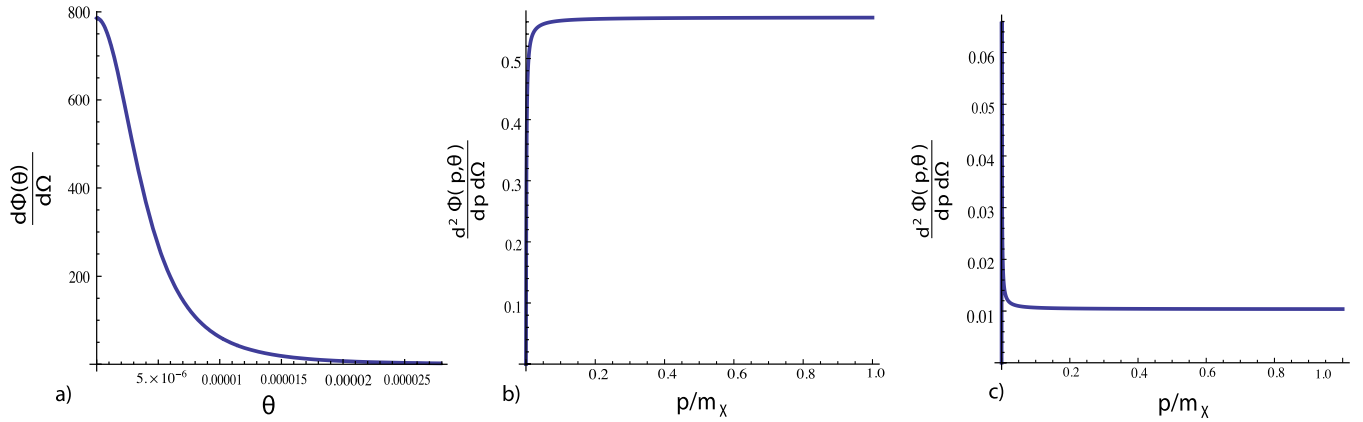


FIG. 2 (color online). Angular and spectral distributions of the photon flux generated by two-photon decays of the mediators with $\gamma v\tau = 0.5R_\odot$ and $\gamma = 1000$. (a) The normalized angular distribution $\frac{d\Phi(\theta)}{d\Omega}/(C_\odot \text{Br}_\nu \text{Br}_\gamma \cdot 10^{-21} \text{ cm}^{-2})$; and (b) the fully differential flux $\frac{d^2\Phi(p,\theta)}{dpd\Omega}/(C_\odot \text{Br}_\nu \text{Br}_\gamma m_\chi^{-1} \cdot 10^{-21} \text{ cm}^{-2})$ in the direction $\theta = 0$, averaged over a Gaussian profile that mimics an angular resolution of $\Delta\theta = 0.1\theta_\odot$; (c) The same for $\theta = 0.1\theta_\odot$.

expression is valid as long as

$$p \in \left[p_0 \sqrt{\frac{1-\beta}{1+\beta}}, p_0 \sqrt{\frac{1+\beta}{1-\beta}} \right] \quad (12)$$

and is understood to vanish for p outside this range. For $l \sin\theta < R_\odot$, the lower limit is again modified to $p_0 \gamma^{-1} (1 - \beta \sqrt{1 - l^2 \sin^2\theta / R_\odot^2})^{-1}$ due to absorption.

The resulting photon spectrum varies dramatically within a tiny range of θ that for the Lorentz boosts considered here is typically well below the angular sensitivity of any gamma-ray telescope. For a detector with angular resolution poorer than the angular size of the Sun, the integration of (11) over $d\Omega$ produces a flat spectrum as long as $p_0/\gamma \leq p \leq p_0 \sqrt{\frac{1+\beta}{1-\beta}}$. However, some detectors, such as Fermi, have an angular resolution smaller than the solar angular size. For demonstration purposes, we take $\gamma \sim 1000$, and average (11) assuming a Gaussian-profile and an angular resolution $\Delta\theta$ of one-tenth the solar size. Since γ is quite large, almost all the photons that reach the detector come from within an angle of $1/\gamma$ that is comparable to the smallest angular size the detector can resolve. Consequently, the signal that would be seen by the detector represents a (relatively) bright central spot of the Sun with an almost exactly flat spectrum. The photon flux drops rather quickly with angle. If detectable, the photon spectrum at $\theta > \Delta\theta$ is again mostly flat except for a very sharp peak near the low momentum end. We illustrate these results in Fig. 2.

Up to this point we have only discussed the spectrum of photons resulting from the decays of long-lived mediators. A similar analysis for charged particles, such as electrons and positrons, is less straightforward due to the complications of the magnetic fields of the Sun and Earth, as well as the solar wind, and their effects on the propagation, energy degradation, and absorption. A proper calculation of the

angular and energy distributions is beyond the scope of this work, but we wish to give a qualitative discussion of these effects, paying particular attention to the implications regarding a possible local solar system component to the PAMELA signal.

A variety of secluded models may be constructed which can lead to electrons and positrons in the final state. It is important to stress that the effects of the magnetic fields and solar atmosphere on these charged states is quite model-dependent, and primarily sensitive to the lifetime and boost of the mediator produced in the WIMP annihilation as well as the production modes of the final state charged particles. If the mediators have a typical decay length on the order of the solar radius or slightly less, processes such as absorption and reflection have the potential to strongly degrade the overall signal flux, and the strong solar magnetic fields may drastically alter the trajectories of the charged particles. If the mediator escapes and decays well past the solar radius, the effects of the heliosphere and the Earth's magnetosphere can be still be significant, especially for less energetic particles in the tens of GeV range and below (where the bulk of the anomalous PAMELA positrons reside).

We must emphasize that the main feature of this signal would be a high degree of anisotropy in the predicted flux of electrons and positrons, correlated with the center of the Sun, which should make this scenario testable. An important consideration is to what degree this tight correlation may be affected or degraded through the processes discussed above. It is instructive to compare this signal with a possible pulsar component to the rising positron fraction of PAMELA [17], where one of the main signatures may again be an anisotropic signal. While the latter anisotropy may be somewhat difficult to detect, this feature in the solar component discussed here would be more pronounced due to the proximity of the Sun and the underlying

production mechanism. This provides additional motivation for studies of the positional, directional, and temporal features of the existing PAMELA and Fermi data sets.

In order to go further and actually attempt to fit the PAMELA and Fermi data, a better understanding of the predicted energy spectrum would be required, which goes beyond the scope of this paper. Of course, this is also quite model-dependent. For example, a simple underlying two-body decay of the mediator to e^+e^- pairs will result in a very hard spectrum, which may not fit the observed spectrum. However, the spectrum can be softened through cascade decays, or the rescattering of mediators inside the solar interior. It would also be important to understand to what extent the electron/positron energy spectrum is affected or softened by interaction with the solar atmosphere and through propagation from the Sun to the Earth.

C. Other local sources

The Sun represents by far the largest astronomical body in the Solar System, and presumably the most efficient WIMP capturing reservoir. However, the lifetime of the mediators can be such that their decays occur deep inside the solar interior and the decays outside are exponentially suppressed. Capture by planets may then be important, and two cases to consider are WIMPs captured by the Earth and Jupiter. The latter option is of interest because of the possibility of very precise observations by means of atmospheric Čerenkov detectors. The flux of WIMP-generated mediators can be estimated using the basic scaling of (2) suitably rescaling the escape velocity, the mass, and the distance to the Earth. Using the results of [11], and assuming for a moment that V decays happen outside the solar interior, we find

$$\frac{\Phi_J}{\Phi_\odot} \sim \left(\frac{v_J}{v_\odot}\right)^4 \times \frac{M_J}{M_\odot} \times \left(\frac{L_{S-E}}{L_{J-E}}\right)^2 \sim 4 \times 10^{-9}, \quad (13)$$

which represents a strong suppression. However, if the decay length of the mediator is less than 5% of the solar radius, the gamma-ray flux from Jupiter may exceed that from the Sun. In this case the use of atmospheric Čerenkov detectors such as HESS, MAGIC and VERITAS, could set additional bounds on the capture rate by Jupiter, and consequently on the WIMP-nucleus cross section.

Finally, the fluxes from the center of the Earth are difficult to present in the same compact form as (2), primarily because of the likely dependence on the annihilation rate. These fluxes are typically much smaller than those generated by WIMP annihilation in the solar interior except for some special cases with resonant energy loss [11]. Nevertheless, the annihilation of secluded WIMPs inside the Earth does allow a probe of much shorter mediator lifetimes, and offers additional signatures in neutrino detectors. In certain models, the signature involves pair production of muons via the decays of the mediators, a

possibility that can be efficiently explored with neutrino telescopes such as SuperK and IceCube.

D. Observational sensitivity

Some of the primary γ -ray observatories, such as atmospheric Čerenkov detectors are of limited utility in this case, as they are unable to directly observe the Sun. While, as noted above, they can be used to place limits on the flux from Jupiter, a fiducial WIMP-nucleon cross section of 10^{-40} cm^2 would generate a flux that is generally too low to provide competitive sensitivity from this source. Thus, we will focus on water Čerenkov detectors and space-borne γ -ray observatories such as Fermi.

The primary limit we will use here comes from the Milagro water Čerenkov detector [18], which has a wide field of view and an angular resolution of 0.75° , slightly larger than the disc of the Sun. The background arises primarily from the scattering of high-energy cosmic rays, producing γ 's which can arrive at the detector from within the Sun's disc. Nonetheless, the limits obtained by Milagro for monochromatic sources are significant and reach up to 10 TeV. We exhibit the ensuing constraint on a contour plot of the γ -flux in Fig. 3, which for characteristic scattering cross sections, already imposes a significant constraint on the electromagnetic branching fractions. The primary flux limits obtained by Milagro are for monochromatic sources, while a monochromatic decay of the mediator as discussed above leads to a smoothed spectrum. Given that the back-

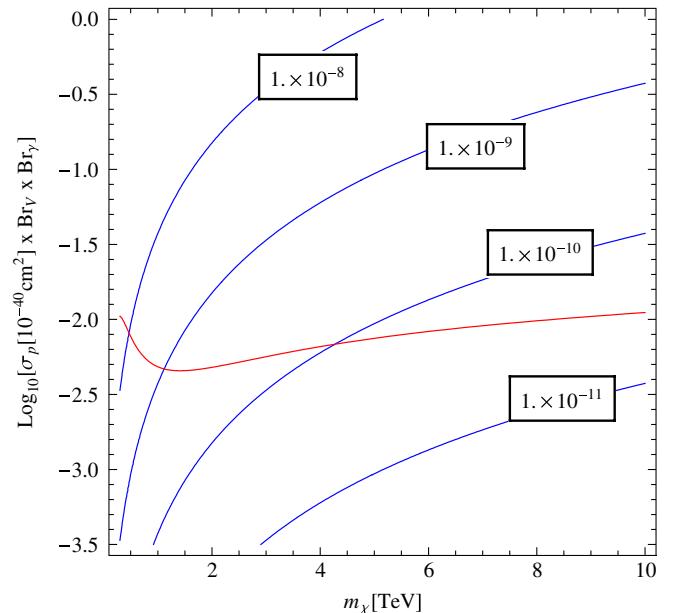


FIG. 3 (color online). Contours (in blue) of the local γ -ray flux Φ in units of $\text{cm}^2 \text{ s}^{-1}$ in the plane of dark matter mass m_χ and a normalized injection cross section assuming spin-dependent scattering. The decay distance in the detector frame has been fixed to R_\odot . The red line indicates a figure of merit for the sensitivity of Milagro (see the text for more details).

ground is a steeply-falling function of energy, while the expected signal is not, we adopt the Milagro bounds on monochromatic sources, weakening it by an order of magnitude in obtaining the bound in Fig. 3.

In the near future, the Fermi-LAT detector, which is also able to observe the Sun [19] with an angular resolution about 10 times better than that of Milagro, should be able to improve on these limits and thus provide a very significant probe of secluded dark matter models. The energy range of Fermi-LAT is also ideally suited to probing secluded WIMP annihilation with masses of a few hundred GeV, where Milagro loses its sensitivity.

It is also possible for long-lived mediators to decay predominantly to electrons and positrons. It is tempting then to consider the implications of these decays in relation to the anomalous positrons fraction observed by PAMELA. As far as we are aware detailed timing and directional analyses have not yet been performed, leaving open the possibility of an intriguing local explanation for these anomalies within the dark matter framework. Given that the existing magnetic fields affect propagation, and the decay chains of the mediators are model-dependent, one could in principle fit the spectral shapes of the observed signals. We have obtained a rough estimate of the required integrated signal flux by using Fermi electron data with an $\sim E^{-3}$ spectrum to infer an estimate of the background and incorporating a new source with a harder power-law spectrum. Depending on the source as well as where the signal turns on (i.e. how large a component of the PAMELA signal) we find a required flux of

$$\Phi_{e^+e^-}(E > 10 \text{ GeV}) \sim 10^{-4} - 10^{-6} \text{ cm}^{-2} \text{ s}^{-1}. \quad (14)$$

We also point out that there will be an $\mathcal{O}(1)$ reduction of the observed signal flux due to the fraction of the time the satellite is facing away from the Sun. Given the assumption that at least part of the excess positron flux does arise from mediator decays, this leads to a *minimal* predicted gamma flux due to the accompanying final state radiation. The typical spectrum of photons resulting from this process would be $(\alpha/\pi)dE/E$, and thus the photon flux will generally be no less than 0.1% of the electron and positron flux, and one infers the following target photon flux:

$$\Phi_{\gamma}^{\text{min}}(E > 10 \text{ GeV}) \gtrsim 10^{-7} - 10^{-9} \text{ cm}^{-2} \text{ s}^{-1}. \quad (15)$$

Thus, while it is possible that the hypothesis of a local origin for the flux anomalies can be directly tested with timing and directional information, it also appears that existing EGRET data [10,20] may already probe the larger range of this associated gamma flux, while it is feasible that Fermi-LAT will be able to cover the entire range of possible fluxes, with a potential $10^{-9} \text{ cm}^{-2} \text{ s}^{-1}$ level sensitivity to multi-GeV gamma rays originating from point sources within a year.

III. SECLUDED MODELS VS γ -RAYS

Having discussed the available sensitivity to γ -rays from the Sun, in this section we will consider some secluded dark matter scenarios which would be subject to this indirect probe. Given that a large variety of model-building possibilities for WIMPs and mediators have been shown to exist [3], we will simply exhibit two classes of models for which the γ -ray flux due to WIMP trapping in the Sun may, for various parameters, be the most sensitive observable and source of constraints. Our approach will be to fix the parameters such that the decay length of the mediator is sufficient to escape the Sun, and then consider whether the scattering cross section leading to capture results in a measurable (or constrained) γ -ray flux according to the limits discussed in the preceding section.

A. Secluded WIMPs with pseudoscalar mediation

A fermionic dark matter candidate χ can be secluded by mediating its interaction with the SM via a pseudoscalar ‘‘axion’’ field a with $m_a < m_\chi$. We will imagine a light WIMP, $m_\chi \sim 10 \text{ GeV}$, as well as a very light mediator, $m_a < 10 \text{ MeV}$. The interactions comprise a series of dimension 5 operators:

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + \bar{\chi}(i\partial_\mu \gamma_\mu - m_\chi)\chi + \frac{1}{2}(\partial_\mu a)^2 - \frac{1}{2}m_a^2 a^2 \\ & + \partial_\mu a \left(\frac{1}{f_\chi} \bar{\chi} \gamma_\mu \gamma_5 \chi + \sum_q \frac{1}{f_q} \bar{q} \gamma_\mu \gamma_5 q \right. \\ & \left. + \sum_l \frac{1}{f_l} \bar{l} \gamma_\mu \gamma_5 l \right) + \frac{\alpha}{4\pi f_\gamma} a F_{\mu\nu} \tilde{F}^{\mu\nu}. \end{aligned} \quad (16)$$

The coupling constants f_i are above the electroweak scale but otherwise are completely arbitrary at this point. Since this is an effective field theory model, depending on the actual UV completion one can achieve the suppression of either f_l^{-1} or f_q^{-1} or both (see, e.g. [21]). The axion can be very long-lived if its mass is low enough to ensure that decays to hadrons and heavier leptons are kinematically forbidden.

We will now briefly discuss the interaction rates pertinent to the γ -ray signal. We note that axion mediation with the WIMP sector has also been explored recently in Ref. [22].

(i) Annihilation

Unless f_χ is parametrically larger than f_f , $\bar{\chi}\chi \rightarrow aa$ will be the dominant annihilation mode given $m_a \ll m_\chi$, and the thermally averaged rate is

$$\langle \sigma v \rangle = \frac{\beta^2}{12\pi} \frac{m_\chi^2}{f_\chi^4} \rightarrow 2.4 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}, \quad (17)$$

where $\beta = (1 - 4m_\chi^2/s)^{1/2}$ is the velocity of the WIMPs in the c.o.m. frame. Note that the annihilation is in the P -wave, a consequence of identical

bosons with overall even parity in the final state. At freeze-out, taking the WIMP velocity to be $\beta^2 \approx 3T_f^2/m_\chi^2 \approx 3/20$, the relic density requires $f_\chi \approx 120 \text{ GeV} \times (m_\chi/10 \text{ GeV})^{1/2}$.

(ii) *Pseudoscalar Decays*

For axions in the MeV range, we can consider decays to photons and electrons,

$$\Gamma_{a \rightarrow \gamma\gamma} = \frac{\alpha^2 m_a^3}{64\pi^3 f_\gamma^2}, \quad \Gamma_{a \rightarrow e^+e^-} = \frac{m_a m_e^2}{2\pi f_e^2}. \quad (18)$$

Depending on the ratios f_e/f_γ and m_a/m_e , either the photon or electron branching may dominate the total width. To maximize the photon fraction, we shall assume that $f_\gamma \sim 10 \text{ TeV}$, while $f_e > \text{few} \times 10^3 \text{ TeV}$, in which case the decays are dominated by the 2-photon final states. Moreover, provided the characteristic coupling of the axions to quarks is small, $f_q > 100 \text{ TeV}$, the axions will be long-lived with a decay length sufficient to escape the Sun,

$$\begin{aligned} L_a &= c\tau_a \gamma_a \\ &\simeq 1.2 \times 10^7 \text{ km} \times \left(\frac{m_\chi}{10 \text{ GeV}}\right) \left(\frac{5 \text{ MeV}}{m_a}\right)^4 \\ &\quad \times \left(\frac{f_\gamma}{10 \text{ TeV}}\right)^2, \end{aligned} \quad (19)$$

where $\gamma_a \simeq m_\chi/m_a \simeq 2 \times 10^3$ for a fiducial normalization of masses and couplings. For the same normalization, one can explicitly check that the absorption cross section of energetic axions in the solar medium is too small to attenuate the flux.

(iii) *Scattering*

The axionlike pseudoscalar mediates spin-dependent WIMP-nucleus elastic scattering, involving the effective axion-nucleon couplings $\mathcal{L}_{a(n,p)} = (\tilde{f}_p^{-1} \bar{p} \gamma^\mu \gamma^5 p + \tilde{f}_n^{-1} \bar{n} \gamma^\mu \gamma^5 n) \partial_\mu a$ where $1/\tilde{f}_{(n,p)} = \sum_q \Delta_q^{(n,p)}/f_q$ in terms of the parameters measuring the spin-content of the nucleons, $\Delta_u^{(p)} = \Delta_d^{(n)} \simeq 0.8$, $\Delta_d^{(p)} = \Delta_u^{(n)} \simeq -0.5$, $\Delta_s^{(p)} = \Delta_s^{(n)} \simeq -0.15$. A straightforward calculation leads to the tree-level cross section, conventionally reexpressed in terms of the ‘‘model independent’’ WIMP-nucleon cross section,

$$\sigma_{p,n} \equiv \frac{\mu_{p,n}^2}{\pi f_\chi^2 \tilde{f}_{(p,n)}^2} \times \begin{cases} 1 & \text{for } m_a^2 \ll 4\mu_N^2 v^2, \\ \frac{16\mu_N^4 v^4}{3m_a^4} & \text{for } m_a^2 \gg 4\mu_N^2 v^2, \end{cases} \quad (20)$$

where μ_N is the reduced mass for the WIMP-nucleus system. Note that depending on the mass of the mediator, WIMP, and type of nucleus involved in the scattering, there is either an enhancement or a suppression. For the trapping rate, we need only

consider WIMP scattering with hydrogen as the scattering is spin-dependent, and in the specific case of the Sun, we must also account for the increase of the characteristic c.o.m. velocity (by a factor of 3–5) due to the WIMPs falling into the Sun’s gravitational well. The characteristic momentum transfer in this case is $|\mathbf{q}| \sim 2m_N v \approx 8 \text{ MeV}$, leading us to consider very light mediators in the MeV range to avoid a possible velocity suppression. A light pseudoscalar with mass in the few MeV range will mediate an enhanced interaction due to the long-range force. Effectively the momentum dependence cancels and we have the usual form for the spin-dependent cross section:

$$\sigma_p \simeq 2.5 \times 10^{-44} \text{ cm}^2 \times \left(\frac{500 \text{ TeV}}{\tilde{f}_p}\right)^2 \left(\frac{10 \text{ GeV}}{m_\chi}\right), \quad (21)$$

where we have used the relic abundance to relate f_χ to m_χ , and we have normalized the effective nucleon couplings to be consistent with the stringent constraints arising from rare K decays, as we will discuss below. We see that direct detection constraints are easily satisfied, as Eq. (21) displays a spin-dependent WIMP-nucleon cross section far smaller than the $\sigma_{(p,n)} < 10^{-37} \text{ cm}^2$ limit [16].

With these results in hand, we see that for light WIMPs and pseudoscalar mediators the solar γ -ray flux can be appreciable. For such light WIMPs in the GeV range, we see from Eq. (2) that the trapping rate can be enhanced by several orders of magnitude compared to weak scale WIMPs, but this is compensated by a generically smaller-spin-dependent cross section shown in Eq. (21).

Interestingly such light mediators are not in conflict with astrophysical bounds when the scale of the interactions f_q is below about 10^6 GeV , because these states will then thermalize in the core of supernovae and so will not be subject to the stringent constraints from cooling. Constraints from BBN cannot rule out one additional thermalized degree of freedom, but in any case a mass in the few MeV range is sufficient to avoid these constraints entirely. Finally, as alluded to above, rare Kaon decays, in particular, searches for $K^+ \rightarrow \pi^+ a$, constrain f_q to be above 100 TeV [21–23], but pose no particular problems for the estimate of the spin-dependent scattering cross section in Eq. (21).

Using Eqs. (2), (3), and (21), and assuming an $\mathcal{O}(1)$ branching of axions to photons and a decay length of order the solar radius as discussed above, we obtain the characteristic gamma-ray flux in the secluded model with axion mediation:

$$\Phi_{\gamma_0} \sim 6 \times 10^{-8} \text{ cm}^{-2} \text{ s}^{-1} \times \left(\frac{10 \text{ GeV}}{m_\chi} \right)^2 \left(\frac{500 \text{ TeV}}{\tilde{f}_p} \right)^2. \quad (22)$$

This is a large flux which is already close to the range probed by satellites like EGRET (although the mass scale is below the sensitivity range of Milagro). We have focused on light WIMPs and mediators, but different parameter ranges may also yield appreciable fluxes at the expense of some fine-tuning of couplings in the quark sector. In particular, a suppression of the coupling to the top quark, with larger light quark couplings, will relax the prohibitive Kaon decay constraints and allow an enhanced capture rate for larger WIMP and mediator masses.

B. Secluded WIMPs with vector mediation

Unlike the axion-mediated models, secluded WIMPs lying in a hidden sector with a spontaneously broken $U(1)_S$ gauge symmetry do not require additional UV completion, and this sector naturally couples to the standard model through the kinetic mixing portal. WIMP scenarios in this framework are straightforwardly formulated [3], and have been a focal point of theoretical interest in the last year due to the positron data released by PAMELA [6]. A general class of WIMP models involve multicomponent states χ , charged under the $U(1)_S$ vector mediator, and the low-energy Lagrangian after symmetry breaking involving the WIMP, vector V , and the Higgs h' , takes the form,

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + \bar{\chi}(iD_\mu \gamma_\mu - m_\chi)\chi + \mathcal{L}_{\Delta m} - \frac{1}{4}V_{\mu\nu}^2 \\ & + \frac{1}{2}m_V^2 V_\mu^2 + \kappa V_\mu \partial_\nu F_{\mu\nu} + \frac{1}{2}(\partial_\mu h')^2 \\ & - \frac{1}{2}m_{h'}^2 (h')^2 + \frac{m_V^2}{v'} h' V_\mu^2 + \dots \end{aligned} \quad (23)$$

If charge-conjugation symmetry is broken via $\mathcal{L}_{\Delta m}$, the Majorana components of the WIMP may be split in mass by $\Delta m \sim \lambda m_V / e'$, which can reduce the elastic scattering cross section and ameliorate constraints on κ from direct detection. The remaining particle physics constraints require that κ be below a few $\times 10^{-3}$.

The lifetime and decay channels for V and h' were analyzed in [24]. There are two regimes in which either V or h' can be very long-lived. The first refers to $m_{h'} > m_V$ and $\kappa < 10^{-9}$. In this case the Higgs' is short-lived, while the vector may have lifetimes in excess of a millisecond. This case is of no interest for us in this paper, because the trapping rate will scale as κ^2 and will be extremely small. The second case with long-lived particles is $m_{h'} < m_V$ and $\kappa \gtrsim 10^{-3}$. In this case the extreme longevity of h' comes from the fact that its decay may only proceed at second order in the mixing angle κ . This kinematic relation renders h' extremely long-lived even for moderately small κ , while the scattering cross section and hence the trapping

rate can remain large. The relevant interaction rates in this case are detailed below.

(i) Annihilation

Once trapped and accumulated in the center of the Sun, WIMP annihilation may lead to various final states: VV , Vh' , and VVV . The latter may only be possible when annihilation proceeds via capture into an $S = 1$ WIMP-onium state [5], but otherwise this annihilation cross section is suppressed by an extra coupling constant. Comparison of the VV and Vh' final state branching is straightforward, once we fix the charge assignment for the Higgs' particle and the spin for the WIMP. If we assume unit charge under $U(1)_S$ for the complex Higgs' field and a fermionic WIMP, then a comparison of the two final states gives:

$$\frac{1}{4} \leq \frac{\langle \sigma_{\chi\chi \rightarrow Vh'v} \rangle}{\langle \sigma_{\chi\chi \rightarrow VV} \rangle} \leq 3, \quad (24)$$

where the brackets include an average over the spin orientation. The upper end of this ratio is achieved when the annihilation proceeds via the formation of WIMP-onium in which case the final state with total spin 1, decaying to Vh' is 3 times more likely than total spin 0, that decays to $2V$. The lower end corresponds to the case when the recombination into WIMP-onium is kinematically forbidden. Notice that this ratio is not changed by Coulomb (Sommerfeld) enhancement of the cross section, as it is identical in both channels. Thus, in this model a minimum of 1 per every 5 annihilation events results in the production of a possibly very long-lived h' particle that is boosted by $m_\chi / m_{h'}$. Note that there is no strict constraint from ensuring the correct relic abundance in this case, as we can take this as a relation which fixes the $U(1)'$ coupling α' , leaving the WIMP mass as a free parameter. For the remainder of this section we will assume the $\chi\chi \rightarrow VV$ mode dominates, in which case $\alpha' \sim 0.02 \times (m_\chi / 500 \text{ GeV})$ [3].

(ii) Higgs' decays

As noted above, the regime of interest here is when $m_{h'} < m_V$ and the Higgs' is long-lived, with the dominant decay $h' \rightarrow l\bar{l}$ occurring at order $\Gamma_h \sim \kappa^4 \times (\text{loop factor})^2$ [24]. Given $m_V \gg m_{h'} \gg 2m_f$ and a boost $\gamma_{h'} \simeq m_\chi / m_{h'}$, the h' decay length is

$$\begin{aligned} L_h = & c\tau_{h'}\gamma_{h'} \\ & \sim 10^7 \text{ km} \\ & \times \begin{cases} \left(\frac{\kappa}{5 \times 10^{-4}}\right)^{-4} \left(\frac{m_{h'}}{500 \text{ MeV}}\right)^{-2} \left(\frac{m_V}{5 \text{ GeV}}\right)^2 & \text{for } m_{h'} > 2m_\mu \\ \left(\frac{\kappa}{5 \times 10^{-3}}\right)^{-4} \left(\frac{m_{h'}}{100 \text{ MeV}}\right)^{-2} \left(\frac{m_V}{500 \text{ MeV}}\right)^2 & \text{for } m_{h'} < 2m_\mu, \end{cases} \end{aligned} \quad (25)$$

where α' has been chosen to fix the relic abundance, and we have assumed a WIMP mass of 500 GeV. However, the longevity of the Higgs' boson does not guarantee its safe passage through the interior of the Sun as there is a potential loss mechanism due to “inverse-Primakoff” type conversion into V on nuclei, $h' + N \rightarrow V + N$, followed by prompt V decay [25]. The cross section for this process on protons can be estimated as follows,

$$\begin{aligned} \sigma_{\text{abs}} &\sim 2\pi\alpha\alpha'\kappa^2 m_V^{-2} \\ &\sim 4 \times 10^{-38} \text{ cm}^2 \times \left(\frac{\kappa}{5 \times 10^{-4}}\right)^2 \left(\frac{1 \text{ GeV}}{m_V}\right)^2. \end{aligned} \quad (26)$$

This $\mathcal{O}(10^{-38} \text{ cm}^2)$ scale for the cross section is sufficiently small that absorption of Higgs' will be negligible, but an increase by 2 orders of magnitude would indeed lead to a significant loss of h' in the solar interior. In particular, a choice of parameters as in the second line of (25) will result in attenuation of the h' flux by more than an order of magnitude. Nonetheless, we see that over much of the parameter space L_h can naturally be large enough for h' 's to escape the Sun without scattering even for relatively large values of κ .

Since the dominant decays are electromagnetic, there will be significant photon production through various processes such as internal bremsstrahlung, etc. However, it is of particular interest to know the branching to γ 's originating from sequential decays of pairs of neutral pions in the product of one-loop induced h' decay, or through a two-loop decay directly to $\gamma\gamma$. The direct decay to $\gamma\gamma$ will have the hardest spectrum, while the π_0 -mediated decay is likely to have larger photonic yield than internal bremsstrahlung. We can estimate these branchings by constructing an effective Lagrangian as follows. We first integrate out all quarks that are heavier than the vector. This leads to an effective interaction of Euler-Heisenberg type $\sim \text{loop} \times (\kappa^2 e^4 / m_Q^4) \times (F_{\mu\nu})^2 (V_{\alpha\beta}^2)$, and similarly a coupling between vectors and gluons. The contribution of such heavy quarks to the relevant decays will thus be suppressed by powers of $(m_V / m_Q)^4$ which we will neglect. We are left with the light quarks u, d, s as well as perhaps c, b if they are lighter than the vector. The next step is to integrate out the vector from the theory at the loop level, which leads to the effective Lagrangian:

$$\mathcal{L}_\psi = -c_\psi \sum_f Q_f^2 \frac{m_f}{v'} h' \bar{\psi} \psi_f, \quad (27)$$

where $c_\psi = -3\kappa^2\alpha/2\pi$, and the sum runs over all quarks lighter than the vector. From here we can straightforwardly compute the decay $h' \rightarrow \gamma\gamma$

much as in the SM, with the result

$$\Gamma_{h' \rightarrow \gamma\gamma} = \frac{\alpha' \alpha^4 \kappa^4}{64\pi^4} \frac{m_{h'}^3}{m_V^2} \left(\sum_f \left| N_c Q_f^4 I\left(\frac{m_{h'}^2}{4m_f^2}\right) \right|^2 \right), \quad (28)$$

where the sum is over all fermions lighter than the vector, and I is the familiar form factor arising from the triangle diagram. The factor in the parentheses is of order one. For a heavy Higgs', $m_{h'} > 2m_\mu$ the dominant decay mode is $h' \rightarrow \mu^+ \mu^-$, and the branching to a pair of photons is quite small, on the order of 10^{-4} . However if h' is lighter than the two-muon threshold, $m_{h'} < 2m_\mu$ the dominant decay mode is $h' \rightarrow e^+ e^-$ with a smaller total width. The branching to two photon pairs in this case can be sizable:

$$\text{Br}_{h' \rightarrow \gamma\gamma} \simeq 10^{-2} \times \left(\frac{m_{h'}}{100 \text{ MeV}}\right)^2 \quad \text{for } m_{h'} < 2m_\mu. \quad (29)$$

For a heavy h' it is still possible to get a large source of gammas through intermediate decays $h' \rightarrow 2\pi^0$, followed by the π^0 fragmenting to photons. The calculation parallels the decay of a light SM Higgs to pions, as detailed in Refs. [26,27]. We consider the theory at even lower energies, below the charm mass, where we can write the effective Lagrangian in terms of the trace of the QCD energy momentum tensor and match onto a chiral Lagrangian. It is then straightforward to calculate the partial width for $h' \rightarrow 2\pi^0$:

$$\Gamma_{h' \rightarrow \pi^0 \pi^0} = \frac{\alpha' \alpha^2 \kappa^4}{2^3 3^4 \pi^3} \frac{m_{h'}^3}{m_V^2} \left(1 - \frac{4m_\pi^2}{m_{h'}^2}\right)^{1/2} |G(m_{h'})|^2, \quad (30)$$

where the function G is

$$G(m_{h'}) \equiv \left\{ \left(\sum_F Q_F^2\right) + \frac{m_\pi^2}{m_{h'}^2} \left[\left(\sum_F Q_F^2\right) + \frac{3}{2} \frac{4z + 1}{1 + z} \right] \right\}, \quad (31)$$

with $z \equiv m_u / m_d \sim 0.56$ and the sum is over the charm and bottom quarks if they are lighter than the vector. This function is numerically $\mathcal{O}(1)$ for the parameters of interest here. The $2\pi^0$ partial width is somewhat smaller than the $h' \rightarrow \mu\mu$ mode, with a branching of

$$\text{Br}_{h' \rightarrow \pi^0 \pi^0} \simeq 5 \times 10^{-2} \times \left(\frac{m_{h'}}{500 \text{ MeV}}\right)^2. \quad (32)$$

Thus approximately 5% of all Higgs' decays will result in 4 photons when $m_{h'}$ is somewhat larger than $2m_\pi$. We conclude that both for a light or heavy h' it

is possible to obtain percent level branchings into photons.

(iii) *Scattering*

The scattering of WIMPs with nuclei in the secluded vector model was discussed in detail in [28], and the regime of most interest here is when a small splitting Δm allows for first-order inelastic scattering on heavier elements. The interesting feature here is that the larger gravitational potential well of the Sun boosts the c.o.m. kinetic energy of the WIMP-nucleus system relative to the case for terrestrial scattering. This leads to a window for Δm which maximizes the trapping rate in the Sun, but which is outside the kinematic range for terrestrial direct detection. Given $E_{\text{kin}}\mu_N$, $\Delta m\mu_N \ll m_V^2$:

$$\sigma_{\text{inel}} = \frac{16\pi Z^2 \alpha \alpha' \kappa^2 \mu_N^2}{m_V^4} \sqrt{1 - \frac{\Delta m}{E_{\text{kin}}}}, \quad (33)$$

which is a simple modification of scattering induced by a finite charge radius for a Dirac WIMP. However, this requires that $\Delta m < E_{\text{kin}} \sim m_\chi^2 v^2 / (2\mu_N)$. The trapping of inelastically scattering WIMPs, first studied in [29] was recently considered in detail in [30,31], and we can directly make use of their results in the present case. Indeed, the trapping in the Sun dominantly occurs through scattering on Fe nuclei, which are lighter than Ge used, for example, in CDMS. However, the larger velocity relevant for trapping in the Sun, means that E_{kin} can be larger allowing for a window in Δm unconstrained by the direct detection limit. Indeed, the cross section above can be large,

$$\sigma_p^{(\text{Fe})} \sim 1 \times 10^{-39} \text{ cm}^2 \times \left(\frac{\kappa}{5 \times 10^{-4}}\right)^2 \left(\frac{5 \text{ GeV}}{m_V}\right)^4 \times \left(\frac{m_\chi}{500 \text{ GeV}}\right), \quad (34)$$

where we have exhibited the equivalent per-nucleon cross section for scattering off Fe nuclei, with α' traded for the WIMP mass through the relic abundance constraint and $m_\chi \gg m_N$. Direct detection constraints then arise from second-order elastic scattering which for large mediator masses in the GeV range are relatively mild, restricting κ to be below 0.1 [28].

An alternative route to maximize the trapping rate while satisfying the direct detection constraints would be to introduce several (two or more) sequentially mixed U(1) groups, all broken at a sub-GeV scale. This would ensure that the WIMP-nucleon cross section contains higher powers of q^2 and the resulting form factor suppresses coherent scattering on nuclei [32]. At the same time, due to the increased

velocity inside the Sun, the degree of suppression in the capture rate will be much smaller.

Given these results, the observable γ -ray flux then follows from the $h' \rightarrow \gamma\gamma$ branching fraction $\text{Br}_\gamma \sim 10^{-2}$ discussed above and the dominant contribution from the WIMP-iron inelastic scattering cross section $\sigma^{(\text{Fe})} \sim 10^{-33} \text{ cm}^2$ to the trapping rate. This easily leads to a detectable flux of gammas for a wide range of parameters,

$$\Phi_{\gamma\odot} \sim 1 \times 10^{-6} \text{ cm}^{-2} \text{ s}^{-1} \times \left(\frac{\kappa}{5 \times 10^{-4}}\right)^2 \left(\frac{5 \text{ GeV}}{m_V}\right)^4. \quad (35)$$

We see that this model leads to a remarkably large gamma-ray flux, that seemingly would have been observed by EGRET and Fermi-LAT. Moreover, for the range $m_{h'} > 2m_\mu$, the accompanying muon flux will generate a flux of ν_μ well in excess of the bounds set by e.g. Super-Kamiokande [33]. However, it is clear that generic choices of parameters may equally well reduce the flux (35) by several orders of magnitude. A more comprehensive scan of the parameter space for this model vs the resulting flux is beyond the scope of this paper, but it clear that the gamma flux does indeed constitute a sensitive probe.

IV. DISCUSSION

Our exploratory study in this paper has suggested a number of striking indirect signatures associated with the annihilation of secluded dark matter trapped in the Sun (and planets), in the form of a high-energy flux of electromagnetic particles tightly correlated with the solar center. In this section, we will finish with a number of additional remarks.

One of the issues that has become apparent is that there are relatively few competitive limits on gamma rays from the Sun, primarily because this is a difficult source to handle for many of the more sophisticated gamma-ray telescopes. However, Fermi-LAT can and has observed the Sun and its impressive angular resolution, and long exposure times point to it as having the best experimental sensitivity for probing the gamma-ray signatures of secluded dark matter scenarios discussed here, with mediator lifetimes on the scale of the solar radius. Indeed, according to our analysis of model scenarios in Sec. III, Fermi-LAT may well provide sensitivity to specific parts of the parameter space which is superior to all other direct or indirect probes.

Beyond direct decays to photons, secluded models also naturally produce a significant branching of mediators to high-energy electron-positron pairs. As discussed earlier, it is an intriguing possibility that massive bodies in the Solar System may provide another possible source of the anomalous electron and positron fluxes in the multi-GeV range. Such scenarios could be probed via high-precision analyses of the spatial or temporal nonuniformities in these

fluxes, particularly the positron fraction observed by PAMELA. Studies of this kind could significantly strengthen the parameter reach in many models of secluded dark matter. Moreover, while the requisite flux of charged particles could be achieved in many variants of secluded dark matter, they most likely will be put to the test by the upcoming Fermi gamma-ray data given the minimal flux that arises through final state radiation.

Within this general framework, we have presented two concrete models of secluded dark matter consistent with the relic abundance and direct detection constraints in which the most accessible experimental signatures are gamma rays and charged particles from the Sun. The crucial ingredient is a mediator with a lifetime long enough to escape from the Sun, but short enough on cosmological or astrophysical scales so that there exist no additional constraints arising from early cosmology beyond those already considered in the literature (see, e.g. [34]). It is likely that other models of this type may be constructed, and indeed it would be worthwhile to explore a more generic scan to understand if the gamma-ray flux observed in these specific models is robust.

Beyond these points that were already touched upon in the text, we would also like to mention some other related issues concerning indirect detection signatures:

- (i) Secluded mediators with long lifetimes can, depending on the particular model and mass range under consideration, greatly enhance the overall neutrino flux reaching the Earth. For example, if the mediators are heavy enough to decay to charged mesons outside the solar radius, the production of muon neutrinos is inevitable. Indeed, even if this decay occurs inside the solar radius, the neutrino yield can be significant provided the decays occur outside the dense core where all muons and pions are quickly absorbed before they are capable of producing energetic neutrinos.
- (ii) There is also the possibility of searching for the annihilation products of standard, neutralino-like WIMPs using gamma-ray telescopes. If the annihilation of WIMPs in the solar interior creates large fluxes of neutrinos, the interaction of neutrinos in the upper layers of the solar atmosphere will result in the production of hadrons (π_0 , K , ...) that decay with the significant yield of γ 's. The surviving γ fraction

which escapes may again be accessible to Fermi-LAT, and can provide alternative sensitivity to the WIMP-powered neutrino flux. While ground-based neutrino telescopes are ideally suited to exploring ν_μ neutrino fluxes, the gamma-ray signature will contain information about other flavors, and be indirectly sensitive to the neutrino energy range where ground-based neutrino telescopes do not have any directional sensitivity.

- (iii) Another possibility that was not considered in this paper is WIMP annihilation to very light quasistable mediators, such as QCD axions. The conversion of axions into photons may occur in the magnetic field of the Sun, resulting in a γ -ray flux. This idea has some similarities with a recent proposal to observe the Sun's transparency to gamma-ray sources that may contain an axionic component [35].

In conclusion, the existence of a generic dark sector with metastable states mediating the interactions between WIMP dark matter and the SM opens up new possibilities for indirect detection signatures in the Solar System. While it is possible to search for charged particles produced by mediator decays, it seems the hard gamma-ray flux produced by decays outside the solar radius may be the most promising "smoking gun" signature due to its tight correlation with the solar center, making it easily distinguishable from cosmic ray-induced backgrounds. As a final remark, we observe that it may be profitable to explore other galactic or extragalactic implications of this scenario in which all stellar objects are effectively imbued with a high-energy gamma spectrum.

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