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Evading direct dark matter detection in Higgs portal models

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ABSTRACT

Many models of Higgs portal Dark Matter (DM) find themselves under pressure from increasingly tight direct detection constraints. In the framework of gauge field DM, we study how such bounds can be relaxed while retaining the thermal WIMP paradigm. When the hidden sector gauge symmetry is broken via the Higgs mechanism, the hidden sector generally contains unstable states which are lighter than dark matter. These states provide DM with an efficient annihilation channel. As a result, the DM relic abundance and the direct detection limits are controlled by different parameters, and the two can easily be reconciled. This simple setup realizes the idea of “secluded” dark matter naturally.

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1. Introduction

The Higgs sector of the Standard Model (SM) enjoys a special feature that it can couple to the hidden sector at the renormalizable level. In particular, a “Higgs portal” interaction term [1–3]

$$V_{\text{portal}} = \lambda_{h\phi} |H|^2 |\phi|^2, \quad (1)$$

where ϕ is a hidden sector scalar, is allowed by all symmetries and has dimension 4. Thus, interactions of this type are expected on general grounds.

An interesting application of this observation is that the Higgs field can couple to dark matter (DM), which is thought to reside in the hidden sector. If the hidden sector is endowed with gauge symmetry, a natural DM candidate would be the corresponding vector field [4–6]. Indeed, the U(1) and SU(N) spontaneous symmetry breaking with a minimal number of scalar fields implies stability of some of the massive gauge fields. This is due to a residual symmetry which acts on the hidden sector states only. In the most general case, multi-scalar systems break CP such that the stabilizing symmetry is Z_2 (or a generalization thereof),

$$A_{\mu}^a \rightarrow (-1)^{n_a} A_{\mu}^a, \quad (2)$$

where a is a group index and n_a is an integer. In the non-Abelian case, this Z_2 is part of a larger unbroken group, e.g. SO(3) or U(1)

[4,7]. It should be noted that for U(1) and SU(2) gauge groups, CP is unbroken since only a single field is required to break the symmetry and the above Z_2 can be viewed as charge conjugation.

The hidden sector “Higgs” field(s) ϕ mixes with the SM Higgs due to the portal coupling Eq. (1). Therefore, the 125 GeV scalar couples to DM, although such a coupling is suppressed by the mixing angle. As a result, the hidden sector DM can annihilate into the SM states and scatter off nuclei as a conventional WIMP would. The current direct DM detection constraints from LUX and PandaX experiments [8,9] are so tight that the WIMP paradigm within the Higgs portal framework finds itself under pressure. The core of the problem is that the couplings controlling DM annihilation and its scattering off nuclei are related, while the direct detection bound requires the latter to be small.

In this work, we emphasize that the Higgs portal models with gauged hidden sectors possess unstable states which can be lighter than dark matter. This provides DM with an efficient annihilation channel which breaks the correlation between the annihilation cross section and the direct detection rate.¹ Thus, all of the constraints can easily be satisfied in this kinematic regime. This type of DM is known as a “secluded WIMP” [11,12], while in this work we show that it is realized quite naturally in the Higgs portal models with gauged hidden sectors. In the context of an SU(2) model, this phenomenon was noted in [4]. The same idea applies of course

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¹ A different mechanism to reconcile relic density constraints with the thermal WIMP paradigm is discussed in Ref. [10] for the scalar Higgs Portal.

to other set-ups as long as there are unstable fields lighter than DM (see e.g. [13,14]).

2. Hidden U(1) sector

The simplest example of a vector DM model with a natural Z_2 symmetry [4,5] (see also [15–17]) is a hidden Abelian gauge sector. Within effective field theory, the model (with a heavy “hidden Higgs”) was analyzed in [18,19].

The Lagrangian is given by

$$\mathcal{L}_{\text{hidden}} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + (D_\mu\phi)^\dagger D^\mu\phi - V(\phi), \quad (3)$$

where ϕ is a charged scalar, $F_{\mu\nu}$ is the U(1) field strength of the gauge field A_μ and $V(\phi)$ is the potential. We take the charge of ϕ to be $+1/2$ for easier comparison to the non-Abelian case. In unitary gauge ϕ can be written as $\phi = (\tilde{v} + \rho)/\sqrt{2}$ where \tilde{v} is the VEV and ρ a real scalar field. The imaginary part of ϕ is eaten by A_μ which obtains the mass $m_A = \tilde{g}\tilde{v}/2$, where \tilde{g} is the gauge coupling. The gauge–scalar interactions are given by

$$\Delta\mathcal{L}_{s-g} = \frac{\tilde{g}^2}{4}\tilde{v}\rho A_\mu A^\mu + \frac{\tilde{g}^2}{8}\rho^2 A_\mu A^\mu. \quad (4)$$

The Z_2 symmetry

$$A_\mu \rightarrow -A_\mu, \quad (5)$$

which is the usual charge conjugation symmetry, makes the massive gauge field stable so that the latter is a viable dark matter candidate.

The visible and hidden sectors interact via the Higgs portal coupling

$$\mathcal{L}_{\text{portal}} = -\lambda_{h\phi}|H|^2|\phi|^2. \quad (6)$$

This coupling leads to the mixing of ρ with the Higgs, which in unitary gauge can be written as $H^T = (0, v + h)/\sqrt{2}$. The fields ρ and h can be written in terms of mass eigenstates $h_{1,2}$ as

$$\begin{aligned} \rho &= -h_1 \sin\theta + h_2 \cos\theta, \\ h &= h_1 \cos\theta + h_2 \sin\theta, \end{aligned} \quad (7)$$

where θ is the Higgs mixing angle and we identify h_1 with the 125 GeV Higgs.

Here we assume that the tree level kinetic mixing between the hypercharge gauge boson and A_μ is zero. This happens if the corresponding generators are orthogonal in the UV completion. For instance, the observable sector can originate from one E_8 factor of the $E_8 \times E_8$ string theory, while the hidden sector comes from the other [20]. The kinetic mixing is not generated radiatively as long as the interaction between the two sectors is due to the Higgs portal term.

Let us now discuss the main phenomenological features of this scenario. All the relevant scattering processes, including DM annihilation and DM scattering on nucleons, proceed through h_1 and h_2 exchange. The dark matter–nucleon interaction cross section is given by (see e.g. [6])

$$\sigma_{A-N}^{\text{SI}} = \frac{g^2\tilde{g}^2}{16\pi} \frac{m_N^2\mu_{AN}^2 f_N^2}{m_W^2} \frac{(m_{h_2}^2 - m_{h_1}^2)^2 \sin^2\theta \cos^2\theta}{m_{h_1}^4 m_{h_2}^4}, \quad (8)$$

where m_N is the nucleon mass, $\mu_{AN} = m_A m_N / (m_A + m_N)$ and $f_N \simeq 0.3$ parametrizes the Higgs–nucleon coupling. One should keep in mind that there is an uncertainty in f_N and here we use the default micrOMEGAS [21] value.

With regard to dark matter annihilation, we focus on the kinematic regime $m_A > m_{h_2}$. For a small $\sin\theta$, the main annihilation

channel in our study is $AA \rightarrow h_2 h_2$ since the relevant vertices are not θ -suppressed. The contributions to this process include the t -channel A -exchange, the s -channel $h_{1,2}$ exchange as well as the contact $AAh_2 h_2$ term. The full cross section expression is quite bulky and not particularly illuminating, thus let us only quote the limit $\sin\theta \ll 1$ and $m_{h_1} \ll m_A, m_{h_2}$,

$$\begin{aligned} \langle\sigma v\rangle &= \frac{\tilde{g}^4}{576\pi m_A^2} \sqrt{1 - \frac{m_{h_2}^2}{m_A^2}} \\ &\times \frac{11m_{h_2}^8 - 80m_{h_2}^6 m_A^2 + 240m_{h_2}^4 m_A^4 - 320m_{h_2}^2 m_A^6 + 176m_A^8}{(4m_A^2 - m_{h_2}^2)^2 (2m_A^2 - m_{h_2}^2)^2}. \end{aligned} \quad (9)$$

In our numerical studies, however, we use the exact result.

In Fig. 1, we display the results of our numerical studies using the package Micromegas [21].² Apart from the direct detection and relic density constraints, we show the perturbativity bounds for the gauge and scalar couplings. In the upper panels, the Higgs decay constraint is also displayed. It comes from the requirement that the LHC Higgs signal strength μ be close to the corresponding SM prediction. When the Higgs production is approximately SM-like, as is the case for small $\sin\theta$, the experimental result $\mu = 1.09_{-0.10}^{+0.11}$ [22] translates into $\text{BR}(h \rightarrow \text{invisible}) \leq 0.11$ at 95% CL. This constrains the h_1 coupling to DM and to h_2 when the corresponding decay channels are open. For the parameter choices of the lower panels, this bound is not relevant. Instead, we display there an upper bound on $\sin\theta$ from the Higgs coupling measurements. The mixing between the SM Higgs and the hidden Higgs suppresses all the couplings of the 125 GeV scalar universally, and is thus subject to the strong LHC bounds.

The thermal WIMP paradigm is consistent with the LUX constraint for $m_A > m_{h_2}$ (plus the resonance regions $m_A \simeq m_{h_{1,2}}/2$) and a sufficiently small θ . If h_2 is relatively heavy, $m_{h_2} \sim \mathcal{O}(100)$ GeV, the required mixing angle is about 0.1 or less. For a light h_2 , the direct detection rate gets quite high and a smaller $\theta \sim \mathcal{O}(10^{-2})$ is necessary. Note also that in the region $m_{h_1} \sim m_{h_2}$, there are substantial cancellations in σ_{A-N} .

The lower panels show that the annihilation cross section at $m_A > m_{h_2}$, m_{h_1} is largely independent of $\sin\theta$. This is because, neglecting the $h_1 - h_2$ mass difference, the gauge boson interactions with $h_{1,2}$ are equivalent to those with ρ (Eq. (4)) which are independent of θ . The s -channel diagrams with $h_{1,2}$ -exchange include θ -dependent Higgs vertices. These however are subleading such that the full cross section depends on θ very weakly. We see from Fig. 1 that the upper bound on $\sin\theta$ from LUX is about 0.1 for the chosen m_{h_2} . Our dark matter candidate evades XENON1T detection for somewhat smaller $\sin\theta$ values, between 10^{-2} and 10^{-1} .

These plots illustrate very well our main point: in the regime $m_A > m_{h_2}$, the direct detection rate is almost uncorrelated with the annihilation cross section.

The SU(2) hidden sector case [4,23,24] is very similar to the Abelian case considered here since both symmetries can be broken by a single field.³ For larger groups, the situation is more involved and in the following section we study an SU(3) example.

² In order to be able to apply the standard thermal freeze-out computations, the dark and visible sectors have to be in kinetic equilibrium at least until DM freeze-out at $T \simeq m_A/20$. This is the case if the interaction rate for elastic scattering of DM particles off SM particles is larger than the Hubble rate. We checked that this is indeed satisfied for the parameter ranges which we consider.

³ Note that no kinetic mixing terms are allowed in case of non-Abelian hidden gauge groups.

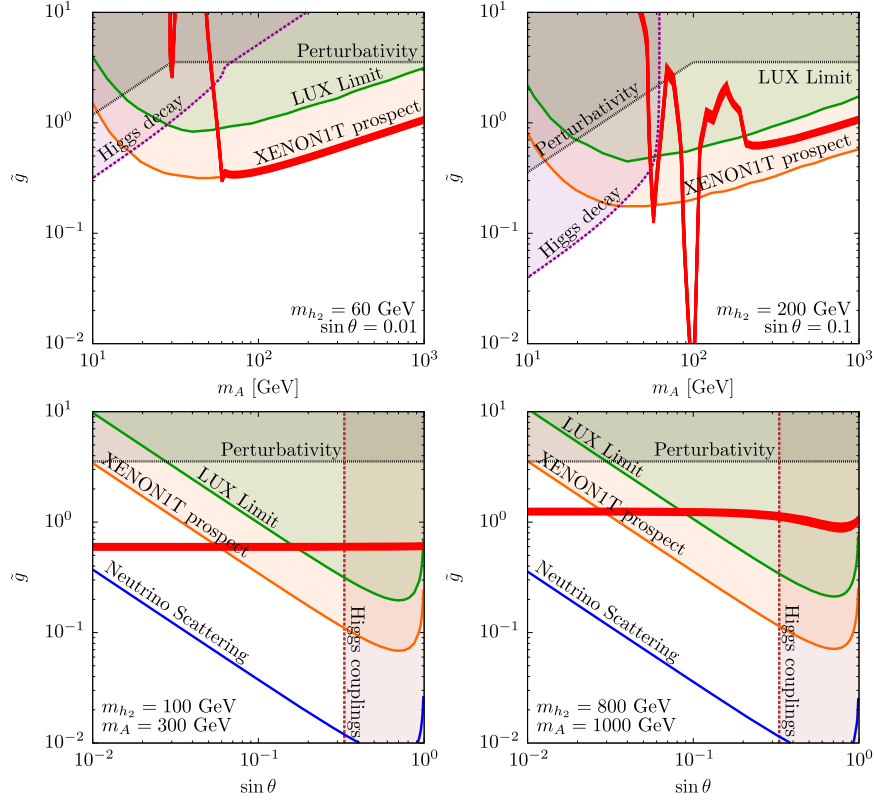


Fig. 1. Dark matter constraints in the plane (m_A, \tilde{g}) (upper panels) and $(\sin\theta, \tilde{g})$ (lower panels) for U(1) DM. The red band indicates the correct relic DM density. The other curves mark the following constraints: grey – perturbativity, purple – invisible Higgs decay, dark red – Higgs couplings, green – LUX 2016 direct DM detection, orange – XENONIT direct DM detection prospects for $2 \cdot y$ exposure. The blue line represents the direct detection event rate corresponding to the relic neutrino scattering off nuclei. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

3. Hidden SU(3) sector

For SU(3) and larger groups [6], the dark matter composition depends on whether CP is broken in the scalar sector. In general, the scalar potential with multiple fields allows for complex couplings which violate CP. In what follows, we will adhere to this general situation, in which case DM is composed of gauge fields.⁴ The corresponding stabilizing symmetry is a global subgroup of the gauge group.

Let us now consider in detail a hidden sector endowed with SU(3) gauge symmetry following Ref. [6]. The symmetry is broken spontaneously (to nothing) by VEVs of two triplets ϕ_1 and ϕ_2 . This is the minimal setup that allows one to make all the SU(3) gauge fields massive. A variation of this model has been considered in [25].

The Lagrangian of the model is

$$\begin{aligned}
 & -\mathcal{L}_{\text{portal}} \\
 & = \lambda_{H11} |H|^2 |\phi_1|^2 + \lambda_{H22} |H|^2 |\phi_2|^2 - (\lambda_{H12} |H|^2 \phi_1^\dagger \phi_2 + \text{h.c.}), \quad (10a)
 \end{aligned}$$

$$\begin{aligned}
 & \mathcal{L}_{\text{hidden}} \\
 & = -\frac{1}{2} \text{tr}\{G_{\mu\nu} G^{\mu\nu}\} + |D_\mu \phi_1|^2 + |D_\mu \phi_2|^2 - V_{\text{hidden}}, \quad (10b)
 \end{aligned}$$

where $G_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu + i\tilde{g}[A_\mu, A_\nu]$ is the field strength of the SU(3) gauge fields A_μ^a , \tilde{g} is the gauge coupling, $D_\mu \phi_i = \partial_\mu \phi_i + i\tilde{g} A_\mu \phi_i$ is the covariant derivative of ϕ_i and H is the Higgs doublet,

which in the unitary gauge can be written as $H^T = (0, v + h)/\sqrt{2}$. The most general renormalizable hidden sector potential can be written as

$$\begin{aligned}
 & V_{\text{hidden}}(\phi_1, \phi_2) \\
 & = m_{11}^2 |\phi_1|^2 + m_{22}^2 |\phi_2|^2 - (m_{12}^2 \phi_1^\dagger \phi_2 + \text{h.c.}) \\
 & + \frac{\lambda_1}{2} |\phi_1|^4 + \frac{\lambda_2}{2} |\phi_2|^4 + \lambda_3 |\phi_1|^2 |\phi_2|^2 + \lambda_4 |\phi_1^\dagger \phi_2|^2 \\
 & + \left[\frac{\lambda_5}{2} (\phi_1^\dagger \phi_2)^2 + \lambda_6 |\phi_1|^2 (\phi_1^\dagger \phi_2) + \lambda_7 |\phi_2|^2 (\phi_1^\dagger \phi_2) + \text{h.c.} \right]. \quad (11)
 \end{aligned}$$

In the unitary gauge, the fields ϕ_1 and ϕ_2 which are responsible for spontaneous SU(3) breaking (to nothing), can be written as

$$\phi_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ v_1 + \phi_1 \end{pmatrix}, \quad \phi_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_2 + \phi_2 \\ v_3 + \phi_3 + iv_4 + i\phi_4 \end{pmatrix}, \quad (12)$$

where the v_i are VEVs and ϕ_i are real scalar fields. In general, CP is broken in the scalar sector and all of the scalar fields mix. In what follows, we make a simplifying assumption that v_3 and v_4 are small and can be neglected. This makes the analysis more tractable without affecting the essence of the model. For definiteness, we also take $v_1 > v_2$.

Our setup enjoys a symmetry that makes the fields A_μ^1 and A_μ^2 stable. To see this, it is sufficient to realize that the model is symmetric under $A_\mu^{1,2} \rightarrow -A_\mu^{1,2}$. This reflection symmetry is part of an

⁴ If CP is preserved, DM may have a pseudoscalar component [7].

unbroken global U(1) which corresponds to the SU(3) gauge transformation $UA_\mu U^\dagger$ with

$$U = e^{i\xi/3} \text{diag}(e^{-i\xi}, 1, 1). \quad (13)$$

Under this symmetry, the gauge field components are rotated as $A^{1(2)} \rightarrow A^{1(2)} \cos \xi \mp A^{2(1)} \sin \xi$ and $A^{4(5)} \rightarrow A^{4(5)} \cos \xi \mp A^{5(4)} \sin \xi$. Thus only $A^{1,2,4,5}$ have non-trivial U(1) quantum numbers, while $A^{3,6,7,8}$ are neutral and can decay into SM matter. The scalar sector has an independent global U(1)' symmetry $\phi_{1,2} \rightarrow e^{i\beta} \phi_{1,2}$. Since U acts effectively as an overall phase transformation on the scalar fields Eq. (12), the vacuum preserves a combination of U(1)' and U . This symmetry has, in particular, the consequence that $m_{A_1} = m_{A_2}$ and $m_{A_4} = m_{A_5}$ (see [6]). It is intact as long as SU(3) is broken in the minimal fashion, that is, via VEVs of only 2 triplets.

Details of the particle spectrum can be found in [6,7], while for our purposes it is sufficient to highlight the following features. In the limit $v_3, v_4 \ll v_1, v_2$ the vector sector is composed of 6 pure states which form 3 mass degenerate pairs with masses

$$m_{A_1}^2 = m_{A_2}^2 = \frac{\tilde{g}^2}{4} v_2^2, \quad m_{A_4}^2 = m_{A_5}^2 = \frac{\tilde{g}^2}{4} v_1^2, \\ m_{A_6}^2 = m_{A_7}^2 = \frac{\tilde{g}^2}{4} (v_1^2 + v_2^2), \quad (14)$$

and two mixed eigenstates

$$A_\mu^{3'} = A_\mu^3 \cos \alpha + A_\mu^8 \sin \alpha, \\ A_\mu^{8'} = A_\mu^8 \cos \alpha - A_\mu^3 \sin \alpha, \quad (15)$$

where

$$\alpha = \begin{cases} \frac{1}{2} \arctan\left(\frac{\sqrt{3}v_2^2}{2v_1^2 - v_2^2}\right) & \text{for } v_2^2 \leq 2v_1^2 \\ \frac{1}{2} \arctan\left(\frac{\sqrt{3}v_2^2}{2v_1^2 - v_2^2}\right) + \frac{\pi}{2} & \text{for } v_2^2 > 2v_1^2 \end{cases} \quad (16)$$

so that $\alpha \in (0^\circ, 60^\circ)$. Their masses are

$$m_{A_3'}^2 = \frac{\tilde{g}^2 v_2^2}{4} \left(1 - \frac{\tan \alpha}{\sqrt{3}}\right), \quad m_{A_8'}^2 = \frac{\tilde{g}^2 v_1^2}{3} \frac{1}{1 - \frac{\tan \alpha}{\sqrt{3}}}. \quad (17)$$

Since $\tan \alpha > 0$, an important consequence of this formula is that

$$m_{A^{1,2}} > m_{A^{3'}}. \quad (18)$$

The lightest fields with non-trivial U(1) quantum numbers are $A_\mu^{1,2}$. They are stable and thus can play the role of dark matter. From now on, we will denote them by A for brevity. Other fields decay into either these fields plus SM states or entirely into the SM final states.

Due to Eq. (18), the DM annihilation channel

$$AA \rightarrow A^{3'} A^{3'} \quad (19)$$

is *always* open and does not suffer the $\sin \theta$ suppression.⁵ $A^{3'}$ is invariant under the transformation of Eq. (13) and thus it decays into SM fields via off-shell scalars. Therefore the “secluded” DM scenario [11] is realized here naturally. The relevant interactions are

$$\mathcal{L} = \frac{\tilde{g}^2}{4} v_2 (-s_\theta h_1 + c_\theta h_2) \\ \times \left[\sum_{a=1,2} A_\mu^a A^{a\mu} + \left(\cos \alpha - \frac{\sin \alpha}{\sqrt{3}} \right)^2 A_\mu^{3'} A^{3'\mu} \right]$$

⁵ In practice, $A^{3'}$ is slightly lighter than $A^{1,2}$ such that this channel incurs some phase space suppression, yet remains efficient.

$$+ \tilde{g} \cos \alpha \sum_{a,b,c=1,2,3'} \epsilon_{abc} \partial_\mu A_\nu^a A^{b\mu} A^{c\nu} \\ - \frac{\tilde{g}^2}{2} \cos^2 \alpha \sum_{a=1,2} \left(A_\mu^a A^{a\mu} \left(A_\nu^{3'} A^{3'\nu} \right) - \left(A_\mu^a A^{3'\mu} \right)^2 \right), \quad (20)$$

where the ϵ_{abc} tensor is antisymmetric in indices 1, 2, 3'. In this expression, we have neglected contributions of heavier scalar and vector states (see the spectrum in [7]).

The analysis of the scalar sector is facilitated assuming small CP breaking. In that case, one can repeat the analysis of [7] while keeping in mind that all the scalars mix and therefore are unstable. Following [7], the lightest spin-0 state can be a “mostly pseudoscalar” χ closely related to φ_4 . Depending on the parameter region, efficient DM annihilation channels

$$AA \rightarrow \chi\chi, \quad h_2 h_2 \quad (21)$$

can be available. These are unsuppressed by $\sin \theta$ and provide a further mechanism to “seclude” [11] our dark matter. The corresponding analytical expressions are bulky and we omit them in this work.

Our numerical results are shown in Fig. 2. At low m_A , the dominant annihilation channel is $AA \rightarrow A^{3'} A^{3'}$ since $A^{3'}$ is always lighter than A , while for heavier DM, the $h_2 h_2$ and $\chi\chi$ final states become important. It is noteworthy that m_A can be as low as 22 GeV at $\sin \theta = 0.01$ without violating any constraints. This is in contrast to the U(1) case where there is no analog of the process $AA \rightarrow A^{3'} A^{3'}$, which excludes very light DM. For $\sin \theta = 0.1$, the effect of the channel $AA \rightarrow h_2 h_2$ or $AA \rightarrow \chi\chi$ is crucial to evade the LUX constraint so that the lowest allowed m_A is about 200 GeV (except for the resonance regions). The lower panels show that again the upper bound on $\sin \theta$ is of order 10^{-1} for electroweak masses.

This analysis can be repeated for larger SU(N) groups (see [6]). In the minimal setting, the gauge symmetry is broken by VEVs of $N - 1$ scalar N -plets. The lightest stable fields correspond to an SU(2) subgroup which gets broken at the last stage. The analog of $A^{3'}$ is lighter than the analogs of $A^{1,2}$ since it mixes with the other Cartan generators. Thus many features of our analysis generalize to SU(N).

4. Conclusion

In the framework of Higgs portal dark matter, we have studied the interplay between the direct DM detection rates and the DM annihilation cross section. Focusing on spin-1 dark matter, we point out that the framework generally contains unstable states lighter than DM, which have a significant coupling to the latter. This opens up an efficient DM annihilation channel into such states thereby effectively decoupling the annihilation cross section from the direct detection rate. The latter is suppressed as long as the mixing angle between the SM Higgs and the “hidden Higgs” is small. This allows us to circumvent the strong LUX/PandaX constraints while retaining the WIMP nature of our dark matter candidate.

We have illustrated this mechanism with both Abelian and non-Abelian vector DM examples. In the U(1) and SU(2) cases, the “hidden Higgs” must be lighter than DM to allow for efficient annihilation. For SU(3) and larger groups, the spectrum automatically contains light unstable vectors which provide DM with an annihilation channel. This argument assumes that CP symmetry is broken in the scalar sector, which in general is the case.

The required mixing angle θ is of order 10^{-1} for $\mathcal{O}(100)$ GeV DM masses. A somewhat smaller θ (by a factor of a few) would suppress the direct detection rate beyond the reach of XENON1T.

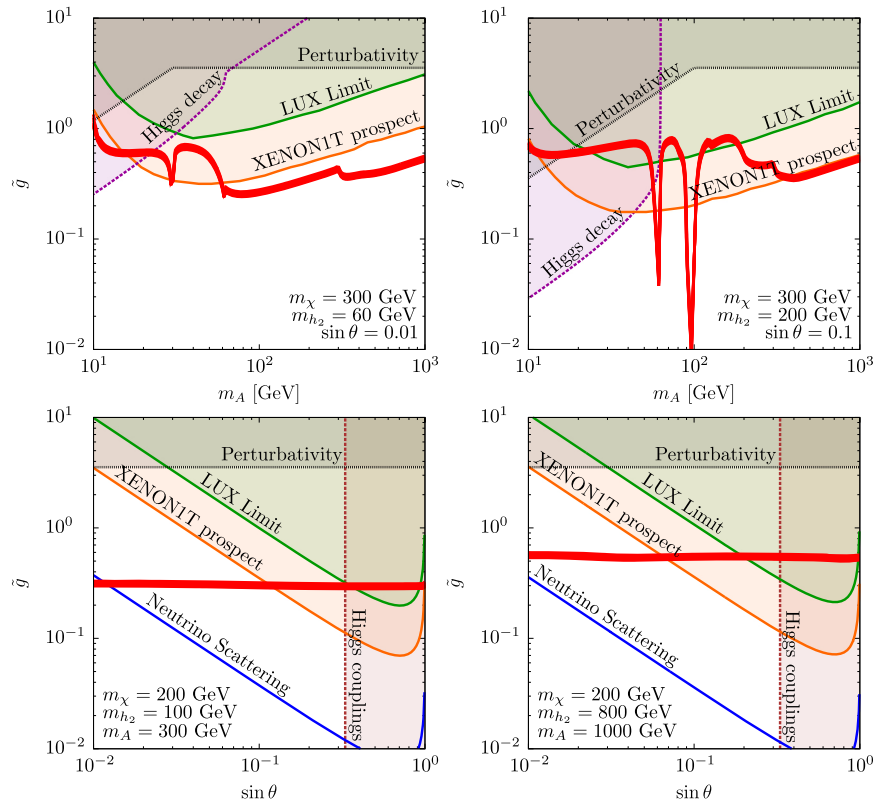


Fig. 2. Dark matter constraints in the plane (m_A, \bar{g}) (upper panels) and $(\sin \theta, \bar{g})$ (lower panels) for SU(3) DM. The red band indicates the correct relic DM density. The other curves mark the following constraints: grey – perturbativity, purple – invisible Higgs decay, dark red – Higgs couplings, green – LUX 2016 direct DM detection, orange – XENONIT direct DM detection prospects for $2 \text{ t} \cdot \text{y}$ exposure. The blue line represents the direct detection event rate corresponding to the relic neutrino scattering off nuclei. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Dark matter can also be quite light, below 100 GeV, in which case the mixing angle is constrained to be in the range $10^{-2} \dots 10^{-1}$, depending on m_A .

We emphasize that the mechanism does not require any significant fine-tuning. When one of the hidden states turns lighter than DM, the annihilation process becomes efficient immediately. The required hidden sector gauge coupling lies in the range 0.1–1 which appears rather natural.

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References

- [1] V. Silveira, A. Zee, *Phys. Lett. B* 161 (1985) 136.
- [2] R.M. Schabinger, J.D. Wells, *Phys. Rev. D* 72 (2005) 093007, arXiv:hep-ph/0509209.
- [3] B. Patt, F. Wilczek, arXiv:hep-ph/0605188.
- [4] T. Hambye, *J. High Energy Phys.* 0901 (2009) 028, arXiv:0811.0172 [hep-ph].
- [5] O. Lebedev, H.M. Lee, Y. Mambrini, *Phys. Lett. B* 707 (2012) 570, arXiv:1111.4482 [hep-ph].
- [6] C. Gross, O. Lebedev, Y. Mambrini, *J. High Energy Phys.* 1508 (2015) 158, arXiv:1505.07480 [hep-ph].
- [7] G. Arcadi, C. Gross, O. Lebedev, Y. Mambrini, S. Pokorski, T. Toma, arXiv:1611.00365 [hep-ph].
- [8] D.S. Akerib, et al., LUX Collaboration, *Phys. Rev. Lett.* 116 (16) (2016) 161301, arXiv:1512.03506 [astro-ph.CO].
- [9] A. Tan, et al., PandaX-II Collaboration, *Phys. Rev. Lett.* 117 (12) (2016) 121303, arXiv:1607.07400 [hep-ex].
- [10] J.A. Casas, D.G. Cerdeno, J.M. Moreno, J. Quilis, arXiv:1701.08134 [hep-ph].
- [11] M. Pospelov, A. Ritz, M.B. Voloshin, *Phys. Lett. B* 662 (2008) 53, arXiv:0711.4866 [hep-ph].
- [12] M. Pospelov, A. Ritz, *Phys. Lett. B* 671 (2009) 391, arXiv:0810.1502 [hep-ph].
- [13] L. Lopez-Honorez, T. Schwetz, J. Zupan, *Phys. Lett. B* 716 (2012) 179, arXiv:1203.2064 [hep-ph].
- [14] M. Duerr, F. Kahlhoefer, K. Schmidt-Hoberg, T. Schwetz, S. Vogl, *J. High Energy Phys.* 1609 (2016) 042, arXiv:1606.07609 [hep-ph].
- [15] Y. Farzan, A.R. Akbarieh, *J. Cosmol. Astropart. Phys.* 1210 (2012) 026, arXiv:1207.4272 [hep-ph].
- [16] S. Baek, P. Ko, W.I. Park, E. Senaha, *J. High Energy Phys.* 1305 (2013) 036, arXiv:1212.2131 [hep-ph].
- [17] M. Duch, B. Grzadkowski, M. McGarrie, *J. High Energy Phys.* 1509 (2015) 162, arXiv:1506.08805 [hep-ph].
- [18] S. Kanemura, S. Matsumoto, T. Nabeshima, N. Okada, *Phys. Rev. D* 82 (2010) 055026, arXiv:1005.5651 [hep-ph].
- [19] A. Djouadi, O. Lebedev, Y. Mambrini, J. Quevillon, *Phys. Lett. B* 709 (2012) 65, arXiv:1112.3299 [hep-ph].
- [20] D.J. Gross, J.A. Harvey, E.J. Martinec, R. Rohm, *Phys. Rev. Lett.* 54 (1985) 502.
- [21] G. Belanger, F. Boudjema, A. Pukhov, A. Semenov, *Comput. Phys. Commun.* 192 (2015) 322, arXiv:1407.6129 [hep-ph].
- [22] G. Aad, et al., ATLAS and CMS Collaborations, *J. High Energy Phys.* 1608 (2016) 045, arXiv:1606.02266 [hep-ex].
- [23] V.V. Khoze, C. McCabe, G. Ro, *J. High Energy Phys.* 1408 (2014) 026, arXiv:1403.4953 [hep-ph].
- [24] A. Karam, K. Tamvakis, *Phys. Rev. D* 92 (7) (2015) 075010, arXiv:1508.03031 [hep-ph].
- [25] A. Karam, K. Tamvakis, *Phys. Rev. D* 94 (5) (2016) 055004, arXiv:1607.01001 [hep-ph].