Indirect and direct signatures of Higgs portal decaying vector dark matter for positron excess in cosmic rays

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Abstract. We investigate the indirect signatures of the Higgs portal $U(1)_X$ vector dark matter (VDM) $X$, from both its pair annihilation and decay. The VDM is stable at renormalizable level by $Z_2$ symmetry, and thermalized by Higgs-portal interactions. It can also decay by some nonrenormalizable operators with very long lifetime at cosmological time scale. If dim-6 operators for VDM decays are suppressed by $10^{16}$ GeV scale, the lifetime of VDM with mass $\sim 2$ TeV is just right for explaining the positron excess in cosmic ray observed by PAMELA and AMS02 Collaborations. The VDM decaying into $\mu^+\mu^-$ can fit the data, evading various constraints on cosmic rays. We give one UV-complete model as an example. This scenario for Higgs portal decaying VDM with mass around $\sim 2$ TeV can be tested by DM direct search at XENON1T, and also at the future colliders by measuring the Higgs self-couplings.

Keywords: dark matter theory, dark matter experiments, cosmic ray experiments

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1 Introduction

There are convincing evidences of nonbaryonic dark matter (DM) in the universe from astrophysical to cosmological scales. According to the results from Planck [1], the dark matter relic density is \( \Omega h^2 = 0.1199 \pm 0.0027 \) with a high precision, while the standard model (SM) for particle physics has no candidate for DM that can account for this measured relic density. We need new physics beyond SM (BSM) for (at least) one new particle playing the role of nonbaryonic dark matter of the universe.

Nonbaryonic dark matter must be stable on cosmological time scale. In case of decay, its lifetime must be much longer than the age of the Universe. The stability of DM is usually guaranteed by imposing a discrete global symmetry, such as \( Z_2 \). There are however some arguments that global symmetries may be generically broken by quantum gravity [2, 3], in which case DM with global charges would be unstable and decay. It can be shown then that the lifetime of DM would be much shorter than the age of the Universe using a naive dimensional analysis, if the DM mass is around electroweak scale, \( \sim O(100) \) GeV–\( O(1) \) TeV (see ref. [4] for example).

Contrary to the global symmetry, a local dark gauge symmetry can be used to guarantee the stability or the longevity of EW scale dark matter. The simplest model would be adding an extra \( U(1)_X \) or some non-Abelian dark gauge symmetry to the SM gauge group \( G_{SM} \) (see refs. [4–8], for example). It is also possible that a hidden sector vector boson could be absolutely stable or its lifetime could be much longer than the age of the universe. Depending on the structure of a given model, the DM could be scalar, fermion or vector particles.
\[ \mathcal{L}_{\text{VDM}} = \mathcal{L}_{\text{ren}}(Z_2 \text{ conserving}) + \mathcal{L}_{\text{non-\text{ren}}}(Z_2 \text{ breaking}) \]

- Thermal relic density
- Direct detection cross section
- Indirect signatures from pair annihilations of VDM’s
- Higgs phenomenology (an additional scalar)

- Suppressed by $1/\Lambda^2$
- $\tau(VDM) \sim 10^{26}$ sec
- Indirect signature from VDM decays: positron excess observed by PAMELA, Fermi and AMS02

Figure 1. Schematic view of the model Lagrangian in this work: the total Lagrangian for the VDM is a sum of the $Z_2$ symmetric renormalizable part and the $Z_2$ breaking nonrenormalizable dim-6 operators, neglecting higher dimensional operators.

One interesting scenario is the so-called Higgs portal Abelian vector dark matter (VDM) model based on $U(1)_X$ dark gauge symmetry (see ref. [9] for example), with an ad hoc $Z_2$ symmetry ($X_\mu \rightarrow -X_\mu$) that stabilizes the VDM. In ref. [21], the authors emphasized that it is important to have a built-in mechanism for generating the VDM mass by introducing a dark Higgs field $\Phi$. The new scalar $\Phi$ would interact with the SM particles due to its mixing with SM Higgs boson through the Higgs portal interaction. There will be two neutral scalar bosons, the mixtures of the SM Higgs boson and the dark Higgs boson. Due to the generic destructive interference between the contributions from two scalar bosons in the amplitude for direct detection cross section, constraints from direct detection experiments such as XENON100, CDMS and LUX can be relaxed significantly and the allowed model parameter space becomes larger than that in the effective model for the Higgs portal VDM [9]. Having a dark Higgs $\Phi$ for the VDM mass, one obtains completely different results compared with the effective VDM model where the VDM mass is given by hand or by Stückelberg mechanism [9].

However, the renormalizable VDM model of ref. [21] may not be the complete theory up to Planck scale, although the model was shown to be perturbative and the electroweak vacuum could be stable up to Planck scale [21]. At some scales, $M_{\text{GUT}} \approx 10^{16}\text{GeV}$ for instance, there could be some new physics which can induce higher dimensional operators for the low energy theory with $G_{\text{SM}} \times U(1)_X$ symmetry, which were not included in ref. [21] (see figure 1). Those nonrenormalizable operators can make the VDM decay after electroweak and dark gauge symmetry breaking, with a resulting DM lifetime that is just at the right order in order to explain the recent observed positron excess [22–24].

In ref. [21], a number of aspects of the renormalizable $U(1)_X$ VDM model have been studied in detail, except for the indirect detection signals. The purpose of this work is twofold. First of all, we work out in detail various indirect signatures and compare with the cosmic ray data, such as $e^+, \bar{p}, \gamma$ or $\nu$ fluxes. There are two different sources of cosmic rays from the VDM. One is the pair annihilations of VDM into the SM particles which are described by $Z_2$ symmetric renormalizable Lagrangian of the VDM model constructed in ref. [21] ($\mathcal{L}_{\text{ren}}$ in figure 1). This part will be constrained by thermal relic density, direct detection and Higgs phenomenology, as described in ref. [21]. The other origin for cosmic rays is the VDM.

\footnote{Extension with non-abelian dark gauge symmetry is also possible to stabilize the VDM [10–15], and fermion DM [16–20].}
decay into the SM particles, which are described by nonrenormalizable higher dimensional operators that break the ad hoc $Z_2$-symmetry ($\mathcal{L}_{\text{non-ren}}$ in figure 1).

In particular we are interested in explaining the positron excess observed by PAMELA, FERMI and AMS02 [22–24], assuming it has the DM-related origin. It turns out that the pair annihilation of (V)DM has difficulties to accommodate the positron excess, because the resulting flux is too small compared with the data without a large boost factor $\sim 10^3$ [28–43]. In general one has to introduce a large boost factor $\sim 10^3$, which however is strongly constrained by the CMB data [44–48] and Fermi/LAT gamma ray measurements [49–55]. Therefore we are led to consider decays of VDM induced by higher dimensional operators.

We write down the complete list of dim-5 and dim-6 operators that cause the VDM decays into various SM particles. Among them, we select operators describing VDM decays into lepton pair $l^+l^-$, and study the positron spectra observed by PAMELA and AMS02. We also present a simple UV completion of the nonrenormalizable operators that could account for the positron data reported by PAMELA and AMS02 Collaborations. In fact a number of works already showed that PAMELA and AMS02 positron excess could be fitted, using $\sim 2$ TeV DM decaying into leptons [56–61]. However thermal relic density or direct detection cross section of the decaying DM for PAMELA and AMS02 were (could) not studied, since these issues are independent of physics for DM decays explaining PAMELA and AMS02.

In this work we fill this gap by assuming that the decaying VDM for positron excess were thermalized by the Higgs portal interaction considered in ref. [21]. If we assume that (i) these positron excess is due to the decaying VDM of mass $\sim 2$ TeV which were thermalized by Higgs portal interaction [21], and (ii) the EW vacuum is stable up to the scale $\Lambda$ where the operators for VDM decays [21], we find that the most parameter region could be probed by the future experiments for direct detection of WIMP’s in the mass range $\sim 2$ TeV. Also the Higgs self-couplings are modified at the level probed at the future colliders such as the ILC. Thus we could make a tight connection between the indirect signature of decaying VDM (with mass $\sim 2$ TeV) from positron excess in cosmic rays and direct detection of such heavy VDM WIMP, as well as the Higgs signal strength and the Higgs-self couplings. These important predictions are newly obtained in this work, compared with other works on decaying DM for positron excess reported by PAMELA, Fermi and AMS02. This accomplishes the second purpose of the present work. Although we work out in the Higgs portal VDM model in this paper, the same strategies could be applied to other types of thermal DMs too.

This paper is organized as follows. In section 2 we give the detailed descriptions of the renormalizable part of the Higgs portal VDM model and then give theoretical and phenomenological constraints on the parameters for a TeV VDM ($X_\mu$) in section 3. In section 4, we show some examples for cosmic ray spectra including gamma ray and neutrinos from the pair annihilation of VDM. In section 5, we list the higher order nonrenormalizable operators relevant for the VDM decays into the SM particles, and present the positron spectra from the VDM decays. We also present one UV-complete model for such a dim-6 operator, as an illustration. Then we show the positron spectra and that $X_\mu \rightarrow \mu^+\mu^-$ could fit the positron spectra for $m_X \sim 2$ TeV, for which we identify the parameter ranges and discuss other observable effects in direct detection of DM and Higgs properties. Finally we give a summary.

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2We assume that the ad hoc $Z_2$ symmetry of the renormalizable Lagrangian for the VDM is accidental symmetry which can be broken by higher dimensional gauge invariant operators, but not by the dim-4 kinetic mixing operators.

3It has to be kept in mind that this excess could be also explained by astrophysical processes [25–27].
2 The Model

We consider a vector dark matter (VDM), $X_\mu$, which is associated with an Abelian dark gauge symmetry $U(1)_X$ implemented with discrete $Z_2$ symmetry $X_\mu \rightarrow -X_\mu$. The simplest renormalizable and unitary model would be the one with an extra complex scalar $\Phi$, whose vacuum expectation value (VEV) is responsible for the mass of $X_\mu$ [21]:

$$L = -\frac{1}{4} X_{\mu\nu} X^{\mu\nu} + (D_{\mu}\Phi)^\dagger (D^\mu\Phi) - \lambda_\phi \left( \Phi^{\dag} \Phi - \frac{v_\phi^2}{2} \right)^2$$

$$-\lambda_{H\Phi} \left( H^\dagger H - \frac{v_H^2}{2} \right) \left( \Phi^{\dag} \Phi - \frac{v_\phi^2}{2} \right) - \lambda_H \left( H^\dagger H - \frac{v_H^2}{2} \right)^2 + L_{SM}. \quad (2.1)$$

Here we neglected the kinetic mixing term $X_{\mu\nu} B^{\mu\nu}$ in order to stabilize the VDM at renormalizable interaction level. The covariant derivative $D_{\mu}$ on $\Phi$ is defined as

$$D_{\mu}\Phi = (\partial_{\mu} + ig_X Q_\phi X_\mu)\Phi,$$

where $Q_\phi$ is the $U(1)_X$ charge of $\Phi$ and it can be rescaled to $|Q_\phi| = 1$.

Assuming the $U(1)_X$-charged complex scalar $\Phi$ breaks $U(1)_X$ spontaneously with a nonzero vacuum expectation value (VEV) $v_\phi$, $\Phi(x) = \frac{1}{\sqrt{2}} (v_\phi + \varphi(x))$, the VDM $X_\mu$ gets mass equal to $M_X = g_X|Q_\phi|v_\phi$, and the hidden sector Higgs field (or dark Higgs field) $\varphi(x)$ will mix with the SM Higgs field $h(x)$ through the Higgs portal interaction, namely the $\lambda_{H\Phi}$ term. The mixing matrix $O$ between the two scalar fields is defined as

$$\begin{pmatrix} h \\ \varphi \end{pmatrix} = \begin{pmatrix} c_\alpha & s_\alpha \\ -s_\alpha & c_\alpha \end{pmatrix} \begin{pmatrix} H_1 \\ H_2 \end{pmatrix} \equiv O \begin{pmatrix} H_1 \\ H_2 \end{pmatrix}, \quad (2.2)$$

where $s_\alpha(c_\alpha) \equiv \sin \alpha(\cos \alpha)$, $H_i(i = 1, 2)$ are the mass eigenstates with masses $M_{H_i}$. $H_1$ will be identifiable as the 125GeV Higgs boson observed at the LHC throughout this paper. The mass matrix of two scalar bosons in the basis $(h, \varphi)$ can be written in terms of either Lagrangian parameters or the physical parameters as follows:

$$M = \begin{pmatrix} 2\lambda_H v_H^2 & \lambda_{H\Phi} v_H v_\phi \\ \lambda_{H\Phi} v_H v_\phi & 2\lambda_\phi v_\phi^2 \end{pmatrix} = \begin{pmatrix} M_{H_1}^2 c_\alpha + M_{H_2}^2 s_\alpha & (M_{H_2}^2 - M_{H_1}^2) s_\alpha c_\alpha \\ (M_{H_2}^2 - M_{H_1}^2) s_\alpha c_\alpha & M_{H_1}^2 s_\alpha + M_{H_2}^2 c_\alpha \end{pmatrix}. \quad (2.3)$$

The mixing angle $\alpha$ of two scalar bosons is determined by

$$\tan 2\alpha = \frac{2M_{12}}{M_{22} - M_{11}}, \text{ or } \sin 2\alpha = \frac{2\lambda_{H\Phi} v_H v_\phi}{M_{H_2}^2 - M_{H_1}^2}. \quad (2.4)$$

This renormalizable Lagrangian for the Higgs portal VDM model, eq. (2.1), has four more parameters compared with the SM: $\lambda_\phi$, $v_\phi$, $g_X$ and $\lambda_{H\Phi}$. For convenience, we shall trade them with the following set of input parameters: $M_X$, $M_{H_2}$, $g_X$ and $\sin \alpha$. Since our aim is to explain the positron excess observed by PAMELA and AMS02 in terms of VDM decays, we will concentrate mainly on heavy VDM with mass around a few TeV in this paper.

Footnote: The issue of $U(1)_X - U(1)_Y$ kinetic mixing is nicely discussed in refs. [62–65].
Figure 2. (Left panel) The horizontal dot-dashed line set the boundary for $\sin \alpha \simeq 0.32$ or $\sin^2 \alpha \simeq 0.1$ from Higgs data. The vertical dot-dashed one marks the limit for perturbativity. The solid(dashed) curves corresponds $\sigma_{XN} = 10^{-44} \left(10^{-45}\right)$ cm$^2$. The vertical red and blue bands set the correct relic density for $M_X = 2$TeV, 3TeV, respectively. (Right panel) This plot shows the relation between $g_X$ and $M_X$ constrained by $\Omega h^2$, the blue band region is allowed with $2\sigma$ variation.

3 $\mathcal{O}$(TeV) VDM and phenomenological constraints

For a successful explanation of the positron excess reported by PAMELA and AMS02, we need a dark matter around $\mathcal{O}$(TeV). Therefore we first would like to show that such a heavy VDM can still be compatible with various constraints from colliders, thermal relic density, theoretical consistencies, etc.. Figures 2, 3 and 5 show that there is indeed an ample parameter space for accommodating a TeV VDM. In this section, we shall provide detailed discussions on various relevant constraints on the renormalizable model Lagrangian eq. (2.1) one by one. The indirect signature from the renormalizable model and from higher dimensional operators will be discussed in section IV and section V, respectively.

3.1 Constraint from Higgs data

The current LHC data on the Higgs signal strengths in various production and decay channels give a constraint on the mixing angle, $|\sin \alpha| \lesssim 0.32$ or $\sin^2 \alpha \lesssim 0.1$ [66]. In the figure 2, the region above the horizontal dot-dashed line yields $\sin \alpha > 0.32$, and therefore is disfavored by the current LHC data. Then assuming the scalar mixing angle $\alpha$ is small, we can make an approximation:

$$\cos \alpha \simeq 1, \quad \sin \alpha \simeq \frac{\lambda_H u_H v_H}{M_{H_2}^2 - M_{H_1}^2}.$$

When $M_X \sim 2$TeV, the relic density constrains the gauge coupling to be around $g_X \sim 0.7$, as shown in figure 2. If we further assume that $H_2$ is still in thermal equilibrium before the VDM freezes out, the $H_2$ mass should be smaller than $M_X$. Taking $M_{H_2} \sim 500$GeV and a
This gives only a rough estimate of approximate values for the parameters. Throughout this section, we restrict the parameters in the following ranges:

\[
0.5 \text{ TeV} < M_X < 3 \text{ TeV},
\]
\[
1 \text{ GeV} < M_{H_2} < 600 \text{ GeV},
\]
\[
0.4 < g_X < 1.0,
\]
\[
0.001 < \sin^2 \alpha < 0.1.
\]

When scanning over these parameters, we take flat distributions in 25 steps for \( M_X, M_{H_2}, \) \( \sin \alpha \) with logarithmic metric, and \( g_X \) with linear metric. The viable and exact values of \( M_X \) and \( g_X \) are further constrained by thermal relic density and perturbativity conditions as discussed below. Distributions of the viable points are illustrated in figures 2, 3, 5 and 6.

The mixing angle \( \alpha \) is also constrained by DM direct search and the lifetime of \( H_2 \) as shown in the right panel of figure 5. The upper bounds are from XENON100 \([67]\) (red), LUX \([68]\) (orange), and vacuum stability (blue) for 2 TeV VDM as examples, where the red and orange regions are excluded. The EW vacuum becomes absolutely stable in the blue region. The lower bound on the mixing angle \( \alpha \) comes from the BBN constraint on the...
lifetime of $H_2$, where we require $H_2$’s lifetime $\tau_{H_2} < 10^{-2}\text{s}$. Otherwise a long-lived $H_2$ could be dangerous to the successful BBN for very small $\sin \alpha$. It turns out that thermalization of the dark sector puts a much more stringent lower bound than BBN except in the low $M_{H_2}$ region. We shall discuss this case later in detail.

### 3.2 Thermal relic density

We are interested in the parameter space, $M_X \sim O(1)\text{ TeV}$ and $M_{H_2} \sim O(100)\text{ GeV}$, aiming at explaining the positron excess. As a result, $\lambda_{\Phi}$ and $\lambda_{H}\Phi$ are small enough that only the first three Feynman diagrams of figure 4 need to be considered for $X \mu$ annihilation. For heavy $X\mu$ and small mixing between $H_1$ and $H_2$, the dominant annihilation channel is $XX \rightarrow H_2H_2$.

Then the quantity $\sigma v$ relevant to thermal relic density is calculated as

$$\sigma v = \frac{1}{3 \times 3 \times 2} \frac{1}{2M_X \sqrt{s}} \int |M|^2 |p_1| d\Omega$$

$$\simeq \frac{g_X^4}{144\pi M_X^2} \left[ 3 - \frac{8 (M_{H_2}^2 - 4M_X^2)}{M_{H_2}^2 - 2M_X^2} + \frac{16 (M_{H_2}^4 - 4M_{H_2}^2 M_X^2 + 6M_X^4)}{(M_{H_2}^2 - 2M_X^2)^2} \right], \quad (3.1)$$

where $\frac{1}{3 \times 3 \times 2}$ accounts for the averaging over polarizations for initial states and identical factor for final states, $s \simeq 4M_X^2$ at decoupling time, and

$$|M|^2 = g_X^4 \left[ 12 - \frac{32 (M_{H_2}^2 - 4M_X^2)}{M_{H_2}^2 - 2M_X^2} + \frac{64 (M_{H_2}^4 - 4M_{H_2}^2 M_X^2 + 6M_X^4)}{(M_{H_2}^2 - 2M_X^2)^2} \right].$$

Since $\sigma v$ is independent of $v$ at the leading order in $v$, we can replace the thermal averaged $\langle \sigma v \rangle$ with eq. (3.1) in the calculation of relic density. For $\langle \sigma v \rangle \sim 3 \times 10^{-26}\text{ cm}^3\text{s}^{-1}$, $M_X \sim \text{TeV}$ and $M_{H_2} \ll M_X$, we have

$$g_X \sim 0.57 \times \left( \frac{M_X}{1\text{ TeV}} \right)^{\frac{1}{2}}. \quad (3.2)$$

As shown in figure 2, the red and blue vertical bands display the correct relic density ($\Omega h^2 = 0.1199 \pm 0.0027$ [1]) of DM for $M_X = 2\text{ TeV}$ and $M_X = 3\text{ TeV}$, respectively. The precise relation between $g_X$ and $M_X$ is shown in the right panel of figure 2, where we used micrOMEGAs3.1 [69] for the numerical calculation.
Figure 5. The left panel shows the scatter plot with direct search constraints from the latest XENON100 [67], LUX [68](red line) and the future XENON1T as function as the dark matter mass and green circles give stable EW vacuum. It can be seen that all green squares can be probed by the XENON1T. The right panel shows the constraints on the mixing angle $\alpha$. The upper bound is from XENON100(red), LUX(orange) and vacuum stability(blue), and the lower bound comes from the BBN constraint on the lifetime and thermalization of $H_2$.

Figure 6. Scatter plots show the Higgs quartic coupling $\lambda_H$ with different x-axis for $M_X \sim O$(TeV). Every point satisfies the relic density in $2\sigma$. Blue triangles are below the LUX Limit and purple circles can be probed by dark matter direct search in the near future. Regions above the horizontal red dotted line give stable EW vacuum.
3.3 Perturbativity and vacuum stability

Perturbativity and vacuum stability of the model can be determined by running RGEs \cite{21} to higher energy scales:

\[
\frac{d\lambda_H}{d\ln \mu} = \frac{1}{16\pi^2} \left[ 24\lambda_H^2 + \lambda_H^2 \Phi - 6g_t^4 + \frac{3}{8} \left( 2g_2^2 + (g_1^2 + g_2^2)^2 \right) - \lambda_H \left( 9g_2^2 + 3g_1^2 - 12g_1^2 \right) \right],
\]

\[
\frac{d\lambda_H \Phi}{d\ln \mu} = \frac{1}{16\pi^2} \left[ 2\lambda_H \Phi (6\lambda_H + 4\lambda_\Phi + 2\lambda_H \Phi) - \lambda_H \Phi \left( \frac{9}{2}g_2^4 + \frac{3}{2}g_1^2 + 6g_1^2 \right) \right],
\]

\[
\frac{d\lambda_\Phi}{d\ln \mu} = \frac{1}{16\pi^2} \left[ 2 \left( \lambda_H^2 + 10\lambda_\Phi^2 + 3g_1^4 \right) - 12\lambda_\Phi g_1^2 \right],
\]

\[
\frac{dg_\chi}{d\ln \mu} = \frac{1}{16\pi^2} \frac{1}{5} g_\chi^4.
\]

For small \( \lambda_\Phi \) and \( \lambda_{H_\Phi} \), the dark sector has negligible effects on the RG running of \( \lambda_H \). Then similarly to the SM, the top quark makes a negative contribution to \( \lambda_H \) from the large top Yukawa coupling \( g_t \), and \( \lambda_H \) would run to a negative value at high scale \( M_\Lambda \), leading to a metastable electroweak vacuum whose lifetime is longer than the age of our Universe. Although the precise \( M_\Lambda \) depends sensitively on \( g_t \) and strong coupling constant \( \alpha_s \), we can use their central values and require positivity of \( \lambda_H \) at scales larger than \( 10^{15} \) GeV. Then we would need \( \lambda_H \gtrsim 0.14 \) at the weak scale, and this would put a constraint on \( M_{H_2} \) and \( \sin \alpha \) from the following relation \cite{21}:

\[
\lambda_H = \frac{M_{H_1}^2 \cos^2 \alpha + M_{H_2}^2 \sin^2 \alpha}{2v_H^2} \gtrsim 0.14.
\]

The allowed parameter space is shown as the blue region in the right panel of figure 5 and the electroweak vacuum is metastable outside of the region. We also show scatter plots for \( \lambda_H \) vs. \( \sin \alpha \) and \( \lambda_H \) vs. \( M_{H_2} \) in figure 6. A sizable deviation from the SM value is possible within the current limits on \( \sin \alpha \) by thermal relic density and direct search for \( X_\mu \). Since the deviation can be as large as \( \mathcal{O}(10\%) \) at tree level, it might be probed at future colliders, such as the ILC for instance. Moreover, all points giving the stable EW vacuum can be tested at XENON1T, as we shall discuss in the following subsection.

The perturbative limit is set by the input value of \( g_X \). We find that \( g_X \lesssim 1.6(1.5) \) can give a perturbative theory up to \( M_{\text{GUT}} \) (Planck scale), respectively. Correspondingly, the VDM mass is bounded from above, \( M_X \approx 7 \) TeV for \( g_X \approx 1.5 \) from eq. (3.2) if nonperturbative effect is neglected.

3.4 Direct search

The VDM \( X_\mu \) can interact with a nucleus through the mixing of \( h \) and \( \varphi \). The cross section of \( X_\mu \)’s scattering off a nucleon is given by

\[
\sigma (X_\mu N \to X_\mu N) = \frac{1}{16\pi} g_X^4 \sin^2 2\alpha \frac{f^2 m_N^2}{v_H^2} \left( \frac{1}{m_{H_2}} - \frac{1}{m_{H_1}} \right)^2 \left( \frac{M_X m_N}{M_X + m_N} \right)^2.
\]

Note that there is a generic cancellation between the \( H_1 \) and \( H_2 \) contributions \cite{21}. When \( M_X \gg m_N \), \( \frac{M_X m_N}{M_X + m_N} \approx m_N \), direct dark matter search experiments will only constrain the product \( g_X^4 \sin^2 2\alpha \), independent of \( M_X \). In figure 2, we show the contours for \( \sigma_{XN} \equiv \frac{1}{16\pi} g_X^4 \sin^2 2\alpha \frac{f^2 m_N^2}{v_H^2} \left( \frac{1}{m_{H_2}} - \frac{1}{m_{H_1}} \right)^2 \left( \frac{M_X m_N}{M_X + m_N} \right)^2 \).
The solid (dashed) curve corresponds \( \sigma_{XN} = 10^{-44} \left( 10^{-45} \right) \text{cm}^2 \), region on the right-handed side gives larger \( \sigma_{XN} \). Note that for large \( M_X \) the XENON100’s bounds \cite{67} are around \( 2 \times \left( \frac{M_X}{1\text{TeV}} \right) \times 10^{-44} \text{cm}^2 \) and LUX \cite{68} improved the limit by a factor of 2.

In figures 3, 5 and 6, we show the scatter plots for the relevant parameters which satisfy the constraints from relic density and LUX and can be probed by the near future XENON1T experiment. We can observe that most parameter space except \( M_{H_2} \approx M_{H_1} \) where the cancellation occurs or no-mixing, \( \sin \alpha \approx 0 \), can be covered by XENON1T. If we require the electroweak vacuum is stable up to high energy scale, then all the allowed points are covered by XENON1T. This is explicitly shown as green squares in figure 3 and green circles in figure 5.

### 3.5 Thermalization of \( X_\mu \) and \( H_2 \)

When calculating thermal relic density of VDM, we are implicitly assuming \( X_\mu \) still has the same temperature as the thermal bath before it freezes out. This is justified as long as \( H_2 \) is in equilibrium\(^5\) since \( X_\mu \) is thermalized by \( X_\mu - H_2 \) interaction. In general, all the relevant processes, such as scattering one \( H_2 + Y \leftrightarrow H_2 + Y \) (\( Y \) is any other particle in the thermal bath), may have to be considered. As an illustration, in the limit of a tiny mixing angle \( \alpha \), we consider one channel for thermalizing \( H_2 \): \( H_2 H_2 \leftrightarrow H_1 H_1 \). This is the most efficient one for thermalization when the temperature is high. We then have the approximate relations:

\[
\Gamma \simeq n_{H_2} \times \langle \sigma v \rangle_{H_2 H_2 \leftrightarrow H_1 H_1}, \quad n_{H_2} \sim T^3, \quad \langle \sigma v \rangle \sim \frac{\lambda_{H\Phi}^2}{T^2}.
\]

In the radiation dominated era, the Hubble constant is \( H \sim T^2/M_{\text{pl}} \), so the condition for equilibrium gives

\[
\Gamma \gtrsim H \Rightarrow \lambda_{H\Phi}^2 \gtrsim \frac{T}{M_{\text{pl}}}.
\]

For \( T \approx 1\text{TeV} \) we have \( |\lambda_{H\Phi}| \gtrsim 10^{-8} \), which in turn leads to \( \sin \alpha \gtrsim 10^{-8} \). In the right panel of figure 5, we show the region under the red dotted line where the dark sector is not thermalized. This constraint is much more stringent than BBN constraint except in the very low \( M_{H_2} \) region. The main purpose of the above discussion is to show that even for very tiny \( \lambda_{H\Phi} \), \( H_2 \) can be still in thermal equilibrium when \( X_\mu \) starts to freeze out.

### 4 Indirect signatures from VDM pair annihilation

Most phenomenology of the VDM with mass \( \mathcal{O}(100\text{GeV}) \) has been studied in ref. [21], except for its indirect signatures. In this section, we focus on indirect signatures from the pair annihilation of VDM (as shown in figure 7) which are described by the renormalizable model Lagrangian, eq. (2.1). Depending on the parameters, the dominant annihilation channels can be different, resulting in different spectra for cosmic rays.

Since it is impractical to show all the cases, here we only discuss 4 different cases, as tabulated in table 1. For a TeV VDM \( (X_\mu) \), the reaction \( X_\mu X_\mu \rightarrow H_2 H_2 \) is the dominant annihilation channel. Therefore the spectrum shape will be truncated at \( M_X \), and fully determined by decay modes of \( H_2 \). The lighter \( X_\mu \) cases, C and D, are chosen just for completeness and comparison with the cases A and B. All these cases are still allowed by

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\(^5\)We require \( H_2 \) is in chemical equilibrium since kinetic decoupling occurs much later and give less stringent constraints.
current experimental constraints considered in the previous section, giving the correct thermal relic density of the VDM, although the dominant annihilation channels for indirect detection for each case could be quite different from each other. Note that the case B would not give an absolute stable vacuum but a metastable vacuum, and here we choose this low $M_{H_2}$ case just for comparison.

The production rate for cosmic rays from DM pair annihilation is given by [72]

$$Q(E, \vec{r}) = \frac{1}{2} \left( \frac{\rho(\vec{r})}{M_{DM}} \right)^2 \sum_i \langle \sigma v \rangle_i \frac{dN_i}{dE}. \tag{4.1}$$

$\frac{dN_i}{dE}$ is the energy spectrum function from a specific annihilation channel $i$, and $M_{DM} = M_X$ in our discussion. The function $\rho(r)$ is the density profile of dark matter. We shall use the Navarro-Frenk-White (NFW) density profile [73],

$$\rho(\vec{r}) = \rho_\odot \left[ \frac{r_\odot}{r} \right] \left[ \frac{1 + (r_\odot/r_c)}{1 + (r/r_c)} \right]^2.$$

Here we use the default values in micrOMEGAs-3.1: $\rho_\odot \simeq 0.3 \text{GeV/cm}^3$, $r_\odot \simeq 8.5 \text{kpc}$ and $r_c \simeq 20 \text{kpc}$ [69]. After production, charged particles propagate through the Galaxy and may lose part of their energy before reaching the solar system. Then the number density $\psi(E, r_\odot)$ can be expressed as [70]

$$\psi(E, \vec{r}_\odot) = \int_{E}^{M_X} dE' \int d^3 \vec{r} G(\vec{r}_\odot, E; \vec{r}, E') Q(E', \vec{r}),$$

where $G(\vec{r}_\odot, E; \vec{r}, E')$ is the Green’s function, paremetrizing the effect during the propagation, such as diffusion and energy loss. Finally the flux is given by

$$\Phi = \frac{v(E)}{4\pi} \psi.$$  

For $\gamma$-ray and neutrinos, after production they travel almost freely, so the fluxes are only dependent on angle region of observation and the integral of squared $\rho$ over the line of sight

$$\Phi_{\gamma/\nu} \propto 2\pi \int \sin \theta d\theta \int_0^\infty dr' \rho^2(r'),$$

where $r' = \sqrt{r^2 + r_\odot^2 - 2rr_\odot \cos \theta}$ and $\theta$ is the angle between the line of sight and the center of Milky Way. We shall use $\theta = \pi/6$ (integrating the region with $\Delta \theta = \pi/60$) as an example.
Table 1. Four cases for illustrating indirect signatures, all are still allowed by current experimental constraints.

| Case | $M_X$ [GeV] | $M_{H_2}$ [GeV] | $g_X$ | $\sin \alpha$ | $(\sigma v)[10^{-26} \text{cm}^3/\text{s}]$ | $\sigma_{XN}[10^{-45}\text{cm}^2]$ |
|------|-------------|----------------|-------|---------------|---------------------------------|----------------------------------|
| A    | 1100        | 280            | 0.56  | 0.08         | 2.25                            | 0.77                             |
| B    | 1100        | 90             | 0.56  | 0.05         | 2.36                            | 0.40                             |
| C    | 400         | 500            | 0.34  | 0.30         | 2.32                            | 2.68                             |
| D    | 400         | 250            | 0.37  | 0.14         | 2.35                            | 0.46                             |

and neglect the $\gamma$-ray induced by inverse Compton scattering and synchrotron radiation from the primary $e^+$ and $\bar{p}$ for simplicity. This is justified as long as we concentrate on the high energy part of the spectrum from $10^{-2}M_X$ to $M_X$. Relative sizes of various contributions are illustrated in [74]. To calculate the cosmic-ray spectra from VDM pair annihilation (i.e. positrons, antiprotons, gammas and neutrinos), we have used micrOMEGAs-3.1 [69] which used Pythia [71] inside.

Generally, the VDM mass determines the energy cut-off of the primary cosmic ray spectra. Since $H_2$ couples to the SM particles with the same pattern as that of $H_1$, its mass determines the branching ratios completely. These different decay final products, together with relevant importance of annihilation channels, can lead different spectra for cosmic rays, as shown in figure 8, although all of them have similar size of $\langle \sigma v \rangle$.

For instance in $e^+$ spectrum, in the case C, the dominant annihilation channel is $X + X \rightarrow W^+ + W^-$, while $X + X \rightarrow H_2 + H_2$ is the dominant one for the case D. About 1/3 of $W$ decay directly to charged leptons and the rest decay hadronically, so that in sum the multiplicity for charged particles is about 20 in a single $W$ decay [75]. While a 250GeV $H_2$ mostly decays to $ZZ$ and $W^+W^-$ whose decay products are then boosted differently. This is the main reason for the different spectra in the case C and D. Similar mechanisms apply to other spectra for $\gamma$, $\bar{p}$ and $\nu's$. For the overall differences between the case A/B and the case C/D, the fluxes at earth have the opposite behavior in some energy ranges, although the primary spectra $dN/dE$ in case A and B are larger than those in case C and D. This is mainly due to the factor $(\rho/M_{DM})^2$ in the source function $Q$, eq. (4.1), and lighter dark matter tends to have larger flux $\Phi$ in the kinematically allowed energy range.

Note that the spectra we discussed so far are only the signatures from dark matter pair annihilation. In reality astrophysical observations of cosmic rays and $\gamma$-ray are the sum of a much larger backgrounds and the above signals. For instance, the positron ($e^+$) flux with $10\text{GeV} < E_k < 300\text{GeV}$ from VDM annihilation in figure 8 has $E^3 \Phi_{e^+}$ around $10^{-6}(\text{cm}^2\text{str s})^{-1}\text{GeV}^2$, while the background $E^3 \Phi_{e^+}$ is about $10^{-3}(\text{cm}^2\text{str s})^{-1}\text{GeV}^2$ which can be inferred from the data in figure 9 of the next section. Therefore, for canonical values of thermal $\langle \sigma v \rangle$, the differences among those cases can hardly be distinguished unless there are some mechanisms for boosting the spectrum, such as Sommerfeld enhancement [76, 77] or Breit-Wigner resonance [78–80] enhancement for explaining the recent observed positron excesses in [22–24]. However, stringent constraints from CMB have been put on such mechanisms for annihilating dark matter [44–48]. Therefore we conclude that the VDM pair annihilations from the renormalizable Lagrangian (2.1) has difficulties to explain the positron excess observed by PAMELA and AMS02. And we shall focus on the decaying dark matter scenario in the next section.

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6 Although possible exceptions exist [81–84], such as multi-component dark matter, halo substructure.
Figure 8. Spectra of $e^+$, $\bar{p}$, $\gamma$ and $\nu$ from vector dark matter annihilation only. The left panels show the primary spectra while the right ones show the spectra at earth after propagation (the units are chosen with the usual convention).
5 Indirect signatures from VDM (X_μ) decay

5.1 Effective operators for decaying VDM (X_μ)

In the renormalizable theory described by the Lagrangian (2.1), the dark matter X_μ can not decay because of the Z_2 symmetry we assumed. This is not true any more if higher dimensional nonrenormalizable operators are taken into account. Generally, higher dimensional operators are suppressed by the power of some new physics scale Λ, above which the nonrenormalizable Lagrangian begins to violate unitarity and fails to describe physical phenomena correctly. In the absence of a complete theory above Λ, we may write down all the operators which are invariant under the gauge group G_{SM} × U(1)_X. To which orders we shall truncate is dependent on the observable we are considering. In the following, we only list higher dimensional operators up to dim-6, especially we focus on those involving both fields from the dark sector and from the SM sector. Such higher dimensional operators for the SM sector upto dim-6 can be found in refs. [85, 86].

Since there are two fields Φ and X_μ in the dark sector, gauge invariant operators in the dark sector would be made of the following operators:

\[ \Phi^\dagger \Phi, \Phi^\dagger \vec{D}_\mu \Phi, X^{\mu\nu}, \tilde{X}^{\mu\nu}, \]

where \( \Phi^\dagger \vec{D}_\mu \Phi = \Phi^\dagger D_\mu \Phi - (D_\mu \Phi)^\dagger \Phi \). The independent effective operators of dim-6 in this sector are

\[ \left( \Phi^\dagger \Phi \right)^3, \left( \Phi^\dagger \Phi \right) \Box \left( \Phi^\dagger \Phi \right), \left( \Phi^\dagger D^\mu \Phi \right)^\dagger \left( \Phi^\dagger D^\mu \Phi \right), \Phi^\dagger \Phi X_{\mu\nu} X^{\mu\nu}, \Phi^\dagger \Phi \tilde{X}_{\mu\nu} X^{\mu\nu}. \]

The operator \((D_\mu \Phi)^\dagger (D_\nu \Phi) X^{\mu\nu}\) is redundant, as it can be shown by partial integration and using equations of motion.

Gauge invariant operators in SM sector are products of the following:

\[ H^\dagger H, H^\dagger \vec{D}_\mu H, B^{\mu\nu}, \vec{B}^{\mu\nu}, \bar{L}_i R_j H, \bar{f}_i \gamma^\mu f_j, (\bar{L}_i \sigma^{\mu\nu} R_j) H, H^\dagger \tau^I H W^I_{\mu\nu}, H^\dagger \tau^I H \tilde{W}^I_{\mu\nu}, \]

where L and R stand for left-handed and right-handed fermion fields, respectively. Note that there is only one dimension-five operator within the SM sector, namely the Weinberg operator for Majorana neutrino masses.

Dimension-6 operators that involve both SM and dark sector fields are

\[ \left( \Phi^\dagger \Phi \right)^2 H^\dagger H, \Phi^\dagger \Phi \left( H^\dagger H \right)^2, \Phi^\dagger \Phi \Box H^\dagger H, \left( \Phi^\dagger \vec{D}_\mu \Phi \right) \left( H^\dagger \vec{D}_\mu H \right), \]

\[ \Phi^\dagger \Phi \left( \bar{L}_i R_j H + h.c. \right), \left( \Phi^\dagger \vec{D}_\mu \Phi \right) \left( \bar{L}_i \gamma^\mu L_j + \bar{R}_i \gamma^\mu R_j \right), \left( \bar{L}_i \sigma^{\mu\nu} R_j \right) H X^{\mu\nu} + h.c., \]

\[ \Phi^\dagger \Phi B_{\mu\nu} X^{\mu\nu}, \Phi^\dagger \Phi \vec{B}_{\mu\nu} X^{\mu\nu}, H^\dagger H B_{\mu\nu} X^{\mu\nu}, H^\dagger H \vec{B}_{\mu\nu} X^{\mu\nu}, H^\dagger H X_{\mu\nu} X^{\mu\nu}, H^\dagger H \tilde{X}_{\mu\nu} X^{\mu\nu}, \]

\[ H^\dagger \tau^I H W^I_{\mu\nu} X^{\mu\nu}, H^\dagger \tau^I H \tilde{W}^I_{\mu\nu} X^{\mu\nu}. \]

The above operators make the whole independent set of operators with both the SM fields and the dark sector fields. Others can be reduced to linear combinations of these operators by using equations of motion.

After the spontaneous gauge symmetry breaking of \( G_{SM} \times U(1)_X \), some of the above effective operators can lead to dark matter \( X_\mu \) decay. Let us consider the following operators
for the VDM decays into two SM particles in the final states:

1. \( \left( \Phi_\mu^i \overrightarrow{D}_\mu \Phi \right) \left( H_\mu^i \overrightarrow{D}_\mu H \right) \Rightarrow X^\mu \to \varphi/h + \gamma/Z, \)
2. \( \left( \Phi_\mu^i \overrightarrow{D}_\mu \Phi \right) \left( \bar{f} \gamma^\mu f \right), \quad \overline{L} \sigma_{\mu\nu} R H X^\mu X^\nu + h.c \Rightarrow X^\mu \to \bar{f} + f, \)
3. \( \Phi_\mu^i \Phi B_{\mu\nu} X^{\mu\nu}, \quad \Phi_\mu^i \bar{B}_{\mu\nu} X^{\mu\nu}, \quad (\Phi \to H) \Rightarrow X^\mu \to \varphi/h + \gamma/Z, \)
4. \( H_\mu^I \tau^I H W^I_{\mu\nu} X^{\mu\nu}, \quad H_\mu^I \tau^I \tilde{H} W^I_{\mu\nu} X^{\mu\nu} \Rightarrow X^\mu \to \varphi/h + \gamma/Z, \)

There are also some interesting three-body decay channels, such as
\( \Phi_\mu^i \Phi \Phi B_{\mu\nu} X^{\mu\nu} \Rightarrow X^\mu \to \varphi + \varphi + \gamma/Z. \)

Generally, three-body decays from these operators are suppressed more compared with two-body decay because of the smaller phase space available. Therefore we will mainly discuss the two-body decay in the following.

5.2 A simple UV completion

It should be pointed out that not all of the above operators need to be investigated simultaneously for the purpose of the positron excess observed by PAMELA and AMS02. The choice is highly dependent on the exact theory beyond energy scale \( \Lambda \) and low energy observables we are interested in. For instance, refs. [87, 88] investigated \( \gamma \)-ray in a similar framework.

Here as a concrete illustration for fermionic final states, let us consider the following operator
\( \left( \Phi_\mu^i \overrightarrow{D}_\mu \Phi \right) \left( \bar{f} \gamma^\mu f \right), \)
which can induce a decay
\( X^\mu \to f \bar{f}. \)

This operator can be induced from the following interactions when both \( \Phi \) and \( f \) are charged under a new extra U(1)‘ symmetry with \( A_\mu' \) gauge field,
\[ \mathcal{L} = (D_\mu' \Phi) \dagger D^{\mu\Phi} + \overline{f} \gamma^\mu D_\mu' f - \frac{1}{4} F'^{\mu\nu} F'^{\mu
u} + (D_\mu' \phi) \dagger D^{\mu\phi} - V \left( \phi^\dagger \phi \right), \]
where the covariant derivatives are
\[
\begin{align*}
D_\mu' \Phi &= (\partial_\mu + ig_X Q_X X_\mu + ig'Q_\Phi A_\mu') \Phi, \\
D_\mu' \phi &= (\partial_\mu + ig'Q_\phi A_\mu') \phi, \\
D_\mu' f &= (D_\mu^{\text{SM}} + ig'Q_f A_\mu') f.
\end{align*}
\]

A new scalar \( \phi \) has been introduced in order to break U(1)‘ spontaneously and make \( A_\mu' \) massive. If only leptons have U(1)‘ charges among the SM particles, then the massive VDM \( X^\mu \) would decay to a lepton pairs only,
\( X^\mu \to l^+ l^-, \quad l = e, \mu, \tau. \)

In such a case, U(1)‘ charge can be identified as lepton number,\(^7\) and \( \phi \) could also couple to right-handed neutrino and give the Majorana mass term after U(1)‘ breaking, acting as

\(^7\)We ignore the anomaly cancellation issue in this paper.
the source of type-I seesaw mechanism. If only $e^\pm$ and $\nu_R$ are U(1)$'$-charged, then $X_\mu$ only decays to $e^+e^-$, see refs. [29, 30] for similar models. For simplicity, we shall assume 100% of $X_\mu$ decay to a single channel for indirect signatures.

In order to explain the positron excess correctly, the lifetime of dark matter should be around $\tau_{DM} \sim 10^{26}\text{s}$, which determines the scale $\Lambda$:

$$\Gamma \sim \left( \frac{g_\Lambda}{\Lambda^4} \right)^\frac{1}{2}, \quad \frac{\tau}{\Gamma} \sim 10^{26}\text{s} \Rightarrow \Gamma \sim 6 \times 10^{-51}\text{GeV}.$$  

For $M = 1\text{TeV}$, we have

$$\Lambda \sim g_\Lambda \left( \frac{M^5\tau}{h} \right)^\frac{1}{4} = g_\Lambda \left( \frac{10^{15}\text{GeV}^5 \times 10^{26}\text{s}}{6.583 \times 10^{-25}\text{GeV} \text{s}} \right)^\frac{1}{4} \sim 2g_\Lambda \times 10^{16}\text{GeV},$$

If $g_\Lambda \sim 0.1$ then $\Lambda \sim 2 \times 10^{15}\text{GeV}$. In the framework of the above U(1)$'$ model, we have the following identifications: $\Lambda \rightarrow M_{A'}$, $g_\Lambda \rightarrow g'$ and $M^5 \rightarrow M_{X}^5 v_F^2$.

Note that those new nonrenormalizable interactions would not affect the VDM annihilation in section III or other results derived from the renormalizable part of the VDM Lagrangian, because the new particles are simply too heavy without doing a precise global fit to the data, focusing only on the channel can or cannot fit the data in a qualitative manner. We shall give simple illustrations without doing a precise global fit to the data, focusing only on the $E_k > 10\text{GeV}$ range.

Since we assume $X_\mu$ can decay to leptons only, it will give rise to indirect signatures in cosmic $e^\pm$, which can be conveniently discussed in terms of two observables: the total flux $\Phi_{e^-+e^+}$ and the positron fraction $\Phi_{e^+}/\Phi_{e^-+e^+}$. Each $\Phi$ is the sum of background flux and the contribution from dark matter decay. The $e^\pm$ background fluxes of interstellar origin can be parametrized analytically as [23, 74]

$$\Phi_{e^-}^{\text{bkg}}(E) = \left( \frac{82.0 \times E^{-0.28}}{1 + 0.224 \times E^{2.93}} \right) \text{GeV}^{-1}\text{m}^{-2}\text{s}^{-1}\text{sr}^{-1},$$

$$\Phi_{e^+}^{\text{bkg}}(E) = \left( \frac{38.4 \times E^{-4.78}}{1 + 0.0002 \times E^{5.63}} + 24.0 \times E^{-3.41} \right) \text{GeV}^{-1}\text{m}^{-2}\text{s}^{-1}\text{sr}^{-1},$$

where $E$ is in GeV unit. For the flux from VDM decay, we calculate it with modifying micrOMEGAs [69]. The production rate is given by

$$Q(E, \vec{r}) = \frac{\rho (\vec{r})}{M_{DM} \tau_{DM}} \frac{dN_{e^\pm}}{dE},$$

(5.1)

d$N_{e^\pm}/dE$ is the energy spectrum function, $M_{DM} = M_X$ in our discussion, $\tau_{DM}$ is the lifetime of $X_\mu$ and $\rho (\vec{r})$ is the density profile of dark matter. We use the NFW profile for the decaying DM, too.

5.3 Decaying VDM ($X_\mu$) and positron excesses

Discussions in the subsection are not entirely new and a number of dedicated model-independent analysis exist in the literature [28, 30, 34–38, 43, 51, 56, 58, 60]. Here we consider on the $X_\mu \rightarrow l^+l^-$, $l = (e, \mu, \tau)$, and shall give a detailed explanation on why each channel can or cannot fit the data in a qualitative manner. We shall give simple illustrations without doing a precise global fit to the data, focusing only on the $E_k > 10\text{GeV}$ range.

For $M = 1\text{TeV}$, we have

$$\Lambda \sim \frac{g_\Lambda (M^5 \Gamma)}{h} = g_\Lambda \left( \frac{10^{15}\text{GeV}^5 \times 10^{26}\text{s}}{6.583 \times 10^{-25}\text{GeV} \text{s}} \right)^{\frac{1}{4}} \sim 2g_\Lambda \times 10^{16}\text{GeV},$$

If $g_\Lambda \sim 0.1$ then $\Lambda \sim 2 \times 10^{15}\text{GeV}$. In the framework of the above U(1)$'$ model, we have the following identifications: $\Lambda \rightarrow M_{A'}$, $g_\Lambda \rightarrow g'$ and $M^5 \rightarrow M_{X}^5 v_F^2$.

Note that those new nonrenormalizable interactions would not affect the VDM annihilation in section III or other results derived from the renormalizable part of the VDM Lagrangian, because the new particles are simply too heavy without doing a precise global fit to the data, focusing only on the $E_k > 10\text{GeV}$ range.

Since we assume $X_\mu$ can decay to leptons only, it will give rise to indirect signatures in cosmic $e^\pm$, which can be conveniently discussed in terms of two observables: the total flux $\Phi_{e^-+e^+}$ and the positron fraction $\Phi_{e^+}/\Phi_{e^-+e^+}$. Each $\Phi$ is the sum of background flux and the contribution from dark matter decay. The $e^\pm$ background fluxes of interstellar origin can be parametrized analytically as [23, 74]

$$\Phi_{e^-}^{\text{bkg}}(E) = \left( \frac{82.0 \times E^{-0.28}}{1 + 0.224 \times E^{2.93}} \right) \text{GeV}^{-1}\text{m}^{-2}\text{s}^{-1}\text{sr}^{-1},$$

$$\Phi_{e^+}^{\text{bkg}}(E) = \left( \frac{38.4 \times E^{-4.78}}{1 + 0.0002 \times E^{5.63}} + 24.0 \times E^{-3.41} \right) \text{GeV}^{-1}\text{m}^{-2}\text{s}^{-1}\text{sr}^{-1},$$

where $E$ is in GeV unit. For the flux from VDM decay, we calculate it with modifying micrOMEGAs [69]. The production rate is given by

$$Q(E, \vec{r}) = \frac{\rho (\vec{r})}{M_{DM} \tau_{DM}} \frac{dN_{e^\pm}}{dE},$$

(5.1)

d$N_{e^\pm}/dE$ is the energy spectrum function, $M_{DM} = M_X$ in our discussion, $\tau_{DM}$ is the lifetime of $X_\mu$ and $\rho (\vec{r})$ is the density profile of dark matter. We use the NFW profile for the decaying DM, too.
Figure 9. These figures show the spectra of the $e^+$ fraction and $e^\pm$ total flux from the decay, $X^\mu \rightarrow l^+l^-$, $l = e, \mu, \tau$ and lifetime $\tau_{DM}$ is chosen to $(4, 2, 0.7) \times 10^{26}$s, respectively. These parameters are chosen for the illustration purpose only. See details in text.

In figure 9, we show the spectra of the positron fraction and $\Phi_{e^+}$ for individual decay channel, $X^\mu \rightarrow l^+l^-$, $l = e, \mu, \tau$. To compare with experimental observation, we have also shown the data from PAMELA, Fermi and AMS02. In the low energy range $E_k < 10$ GeV, it is known that solar wind can have significant effects on the charged particles, the so-called solar modulation which depends strongly on the solar activity. Since the uncertainty for the background flux in this range is large, we shall not discuss the spectra for $E_k < 10$ GeV any further in this paper. For $E_k > 10$ GeV the mass of $X_\mu$ and lifetime $\tau_{DM}$ are chosen to give relatively better fit with the data.

From green dotted curves in figure 9, we can see that $X^\mu \rightarrow e^+e^-$ can not be consistent with the positron fraction and the total flux simultaneously. The reason is that the spectrum of $e^\pm$ from $X_\mu$ is very hard and too sharp around $E = M_X/2$, and it is inconsistent with the Fermi data on the total flux.

In case of $X_\mu \rightarrow \mu^+\mu^-$, the situation is much better as shown in the blue dot-dashed curves for $M_X = 2$ TeV and $\tau_{DM} \approx 2 \times 10^{26}$s. Since the produced $\mu^\pm$ undergoes subsequently three-body decay $\mu^\pm \rightarrow e^\pm + \nu_e + \nu_\mu$, the resulting $e^\pm$ spectrum from the VDM decay becomes much softer compared with the $X^\mu \rightarrow e^+e^-$ case.

In the $\tau^+\tau^-$ case, the $e^\pm$ spectrum is even softer than the $\mu^+\mu^-$, since only one third of $\tau^\pm$ decay to $\mu^\pm$ and $e^\pm$. Other $\tau$’s decay hadronically into lighter mesons which then decay further to pions, followed by $\pi^\pm \rightarrow \mu^\pm + \nu_\mu$ and $\pi^0 \rightarrow 2\gamma$. However, the spectrum’s softness could be compensated with an even heavier $X_\mu$. As we show in the brown double-dot-dashed lines of figure 9, $M_X = 7$ TeV and $\tau_{DM} \approx 0.7 \times 10^{26}$s can give a good fit with the data. We shall note $m_X = 7$ TeV lies in the boundary of previous constraints as the perturbativity limits $g_X \lesssim 1.5$ which further set the upper bound $M_X \lesssim 7$ TeV to give correct relic density.

\textit{\gamma-ray Constraints:} it is well known that the cosmic $\gamma$-ray is an important constraint on both pair-annihilating and decaying dark matter. If combined with gamma-ray constraint for decaying dark matter based on Fermi-LAT data [28, 49, 52-55], the only viable channel is $X_\mu \rightarrow \mu^\pm\mu^-$. The reason is that for the $e^+e^-$ channel all the lost energy goes to photons. For the $X_\mu \rightarrow \tau^+\tau^-$ case, the decay products has a lot of $\pi^0$ which then all decay to $2\gamma$.
while for $X_\mu \rightarrow \mu^+ \mu^-$ a large part of muon energy is carried by the neutrinos in the decay product of muon, leaving less energy for electron to radiate $\gamma$. Interpretation and constraints after AMS02 have been discussed in [32–39, 41–43, 51, 56–61], which would not change the $\gamma$-ray constraints.

5.4 Implications for thermal VDM with mass $\sim 2$ TeV

Accounting for the positron excess observed by PAMELA and AMS02 through thermal VDM ($\sim 2$ TeV) decaying into $\mu^+ \mu^-$ will restrict the parameter space of the renormalizable Lagrangian, which is one of the main results of this paper. As we have shown in section 3, for $\mathcal{O}$(TeV) VDM, the thermal relic density can pin down the gauge coupling $g_X \simeq 0.76$ in the dark sector (see figure 2 and eq. (3.2). Then only $M_{H_2}$ and the exact mixing angle $\alpha$ are not fixed, but they are correlated with and constrained by Higgs data, DM direct searches, BBN and thermalization assumption, displayed in figure 5. Taking $M_{H_2} \simeq 300$ GeV as an example, we have $\sin \alpha \lesssim 0.3$ and $\lambda_H \gtrsim 0.129$.

With all the current constraints taken into account and taking $M_{H_2} \geq 150$ GeV, sizable deviation is possible for the Higgs self-coupling as shown in figure 6. Precise measurements of Higgs self couplings at the future colliders then could fix $M_{H_2}$ and $\alpha$. Then we can predict the $X_\mu$-nucleon cross section for DM direct searches and our model gets testable. If we further require that the electroweak vacuum is stable up to the scale $\Lambda \sim 10^{15} - 10^{16}$ GeV, then all parameter space can be probed by XENON1T. This is an interesting and important result within our approach on decaying VDM thermalized by Higgs portal interaction.

6 Summary

In this paper, we have investigated the phenomenology (mainly focusing on indirect signatures) of a vector dark matter $X_\mu$ in the framework of Higgs portal model, enlarging the SM gauge group $G_{SM}$ by a dark $U(1)_X$. We first discussed the primary cosmic rays, including $\gamma$-ray and neutrino fluxes, from $X_\mu$-$X_\mu$ annihilation and compare the spectra in several cases. In order to explain the positron excess observed by PAMELA and AMS02, we then focus on the TeV scale $M_X$ and show it can evade all the constraints from the Higgs data, relic density, perturbativity and dark matter direct search. Signals from heavy $X_\mu$ pair annihilation into leptons are well below the background and data. Since having the boost factor from the Sommerfeld enhancement is strongly constrained and basically ruled out by CMB, we then turn to the signatures from $X_\mu$’s decay for explanation of the positron excess observed by PAMELA and AMS02.

We have also presented all the independent dim-6 operators that involve both standard model and dark sector particles, and that are invariant under the $G_{SM} \times U(1)_X$ gauge symmetry. After the breaking of $G_{SM} \times U(1)_X$, the VDM $X_\mu$ can decay to the SM particles. A TeV VDM $X_\mu$ can also explain the excess of positron fraction recently observed in PAMELA, FERMI and AMS02 experiments. We give an example model to implement a leptophilic interaction and show the relevant indirect signature. It is shown that $X_\mu \rightarrow e^+ e^-$ gives a spectrum too hard to explain the observation while $X_\mu \rightarrow \mu^+ \mu^-$, $\tau^+ \tau^-$ can be consistent with both positron fraction and the total $e^\pm$ flux. However, if we take the constraints from the gamma ray, then only $X_\mu \rightarrow \mu^+ \mu^-$ is viable.

Our study presented in this paper is different from other model independent analysis of cosmic rays in the literature. We demonstrated explicitly that thermalization of the VDM is possible for $\sim$ TeV scale VDM, and then considered the VDM decays into a lepton pair. The
indirect searches for cosmic rays can determine the VDM mass, which then fixes the $U(1)_X$ gauge coupling for giving the thermal relic density. The only left two correlated parameters are the mass of second scalar and its mixing angle with Higgs. These two can be further probed by future collider searches, for instance, precision measurement of Higgs self coupling or production of the second scalar, and DM direct searches at XENON1T for example. The physical observables we have discussed systematically in this paper are complementary to each other and testable in terrestrial experiments. Similar analyses could be done for other types of decaying DM assuming they are thermalized through some interactions (such as Higgs portal or singlet portal interactions [15]).

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