

Higgs in Space!

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Abstract

We consider the possibility that the Higgs can be produced in dark matter annihilations, appearing as a line in the spectrum of gamma rays at an energy determined by the masses of the WIMP and the Higgs itself. We argue that this phenomenon occurs generally in models in which the dark sector has large couplings to the most massive states of the SM and provide a simple example inspired by the Randall-Sundrum vision of dark matter, whose 4d dual corresponds to electroweak symmetry-breaking by strong dynamics which respect global symmetries that guarantee a stable WIMP. The dark matter is a Dirac fermion that couples to a Z' acting as a portal to the Standard Model through its strong coupling to top quarks. Annihilation into light standard model degrees of freedom is suppressed and generates a feeble continuum spectrum of gamma rays. Loops of top quarks mediate annihilation into γZ , γh , and $\gamma Z'$, providing a forest of lines in the spectrum. Such models can be probed by the Fermi/GLAST satellite and ground-based Air Cherenkov telescopes.

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1 Introduction

The cosmological evidence for dark matter (DM) is overwhelming, and yet we are still unsure of its nature. We approach a new era of searching for dark matter, with the current generation of experiments closing in on weak scale masses and couplings. Among the many alternate possible visions of dark matter, Weakly-Interacting Massive Particles (WIMPs) are among the most appealing, largely because of the potential to understand their population as a thermal relic, and the prospect of a connection to the dynamics of electroweak symmetry breaking (EWSB) [1].

If WIMPs are indeed part of the dynamics of EWSB, they are expected to have relevant interactions with the fields of the Standard Model. The next few years are likely to be very rewarding, as WIMPs may be produced in high energy particle accelerators, scattering non-relativistically with heavy nuclei, and/or annihilating into observable particles in space. Each process represents a unique opportunity to learn about the nature of dark matter. In particular, indirect detection of two WIMPs annihilating in our galaxy into high energy gamma rays could potentially be our first indication that WIMPs are electroweakly active particles, and our first non-gravitational glimpse of dark matter.

WIMPs are dark, which implies that they do not couple directly to photons. The processes by which observable gamma rays are produced are thus typically complicated. In most models, the dominant annihilation is into charged particles which can themselves radiate photons, hadronic states including π^0 's which decay into $\gamma\gamma$, and/or heavier states which decay into quarks and leptons. The continuum of photons thus produced receives some imprint of the WIMP and its annihilation channels, but strong features are often lacking. More striking are the (typically) subdominant $2 \rightarrow 2$ reactions in which a WIMP annihilates directly into a photon and another particle, X . The kinematics produce a line in the photon energy,

$$E_\gamma = M \left(1 - \frac{M_X^2}{4M^2} \right) \quad (1)$$

with M and M_X respectively the WIMP and X masses, which is a striking feature compared to expected astrophysical backgrounds. γ line signals could also be produced by decaying dark matter, see *e.g.* [2]. In this case, Eq. 1 is replaced by $E_\gamma = \frac{M}{2} \left(1 - \frac{M_X^2}{M^2} \right)$. Such emission could be detected by current gamma-ray telescopes, such as Fermi-LAT, CANGAROO, HESS, MAGIC, and VERITAS, which cover energies in the range of GeV - 10 TeV with resolutions of order $\Delta E/E \sim 0.1$. Depending on the WIMP mass and couplings, several particles may play the role of X , thus an entire forest of lines may be produced by WIMP annihilations [3], and their positions and intensities may reveal many features of the underlying theory of dark matter, including the presence of heavy states no longer present in the Universe today.

It was recognized early on that supersymmetric dark matter may produce γZ as well as $\gamma\gamma$ final states [4]. However, the finite energy resolution of the experiments makes identifying the γZ line distinctly from the $\gamma\gamma$ line very challenging unless the WIMP is very light, like in the Inert Higgs DM Model [5]. In theories with more massive states, such as the $B^{(1,1)}$ Kaluza-Klein modes present in the 6d chiral square [6] model of universal extra dimensions [7], the process $B_H B_H \rightarrow \gamma B^{(1,1)}$ may reveal the presence of $B^{(1,1)}$ and help distinguish the chiral square from the 5d case [8]. Theories with light Z 's that couple to WIMPs may also have a strong line from annihilation into γZ [9].

An interesting possibility is when X is the Higgs boson h . The apparent coincidence between the relic density and the electroweak scale raises the possibility that the dynamics of dark matter is related to that of the electroweak symmetry-breaking, and such a connection suggests that WIMPs may have important couplings to massive states, such as top quarks, electroweak bosons, and the Higgs. The observation of a line in the gamma ray spectrum whose position reflects the Higgs mass would be an exciting discovery, and could even constitute the first observation of a Higgs production process. However, it would still require observation and measurement of the Higgs at colliders to decisively identify the particle associated with that line as the Higgs. As Tevatron and LHC experiments strive to produce the Higgs in particle collisions, it is fascinating that dark matter annihilations may already be producing it in space. Identification of a Higgs line would itself tell us much about the theory of WIMPs. It would give credence to the notion that WIMPs are part of the dynamics of electroweak symmetry-breaking, and suggest properties such as the WIMP spin. Scalar WIMPs carry no intrinsic angular momentum, and Majorana statistics allow s -wave annihilations only through a spin-singlet state, either of which lead to suppression

of WIMP annihilations into γh by the tiny WIMP velocities v expected in the galaxy. In principle, vector dark matter could lead to a γh signal. However, in popular theories such as Little Higgs models [10] and Kaluza-Klein dark matter [11] annihilation into γh results from box diagrams and is highly suppressed. The observation of a γh line would thus favor particular combinations of WIMP spins and interactions.

In this article, we explore the possibility that WIMPs may annihilate into γh , producing Higgs bosons in space in association with a photon whose energy reflects the mass of the Higgs (and the WIMP). Our calculations are performed in the context of an effective theory describing a Dirac fermion dark matter candidate. Such a model captures the low energy physics of the Randall-Sundrum (RS) model of a warped extra dimension [12], in which the WIMP is a right-handed neutrino whose stability results from the need to suppress rapid proton decay [13]. However, we emphasize that this model is not unique – as we argue in Section 2, a WIMP whose dynamics is intimately linked with EWSB (and has appropriate statistics) is *likely* to annihilate in this way. In Section 3 we present the continuum of gamma rays expected when WIMPs couple strongly to top quarks, and we compute the expected intensity of the γZ , γh , and $\gamma Z'$ lines and discuss prospects for detection. We briefly discuss predictions for the relic density, direct detection, and some collider signals in Section 4. In Section 5, we finish with an outlook and conclusions.

2 A Dark Matter–Top Quark Connection

If dark matter arises as part of the dynamics of electroweak symmetry-breaking, it is natural to expect the WIMP to have couplings which favor the most massive states of the Standard Model. Here, we explore the possibility that the WIMP has important couplings to the top quark, through which it can couple at the loop level both to photons and to Higgs bosons.

2.1 Effective Theory

We take the WIMP to be a Dirac fermion ν which is a singlet under the SM gauge interactions. It is charged under a (spontaneously broken) $U(1)'$ gauge symmetry, the massive gauge boson of which acts as a portal to the SM by coupling to the top quark. The effective Lagrangian contains,

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{SM} - \frac{1}{4} \hat{F}'_{\mu\nu} \hat{F}'^{\mu\nu} + M_{\hat{Z}'}^2 \hat{Z}'_\mu \hat{Z}'^\mu + \frac{\chi}{2} \hat{F}'_{\mu\nu} \hat{F}'^{\mu\nu}_Y + \hat{g}_t^{Z'} \bar{t} \gamma^\mu P_R \hat{Z}'_\mu t \\ & + i\bar{\nu} \gamma^\mu \left(\partial_\mu - i\hat{g}_\nu^{Z'} P_R \hat{Z}'^\mu \right) \nu + M_\nu \bar{\nu} \nu \end{aligned} \quad (2)$$

where $\hat{F}'_{\mu\nu}$ ($\hat{F}'^Y_{\mu\nu}$) is the usual Abelian field strength for the \hat{Z}' (hypercharge boson), $\hat{g}_t^{Z'}$ is the \hat{Z}' coupling to right-handed top quarks, and $\hat{g}_\nu^{Z'}$ is the coupling to right-handed WIMPs. M_ν is the WIMP mass. One can easily include a coupling to the left-handed top (and bottom). Our choice to ignore such a coupling fits well with typical RS models, balancing the need for a large top Yukawa interaction with control over corrections to precision electroweak observables. The parameter χ encapsulates the strength of kinetic mixing between the Z' and SM hypercharge bosons, and the hatted (unhatted) quantities are those before (after) mixing, as discussed below.

We have included hypercharge- \hat{Z}' kinetic mixing through the term proportional to χ . Such a term is consistent with the gauge symmetries, and even if absent in the UV, will be generated in the IR description by loops of top quarks¹. The kinetic mixing parameter χ generates an effective coupling of SM states to the \hat{Z}' , and through electroweak symmetry breaking, mass mixing of the \hat{Z}' with the SM Z gauge boson resulting in a coupling of ν to the SM Z boson. In terms of the quantities,

$$\eta \equiv \frac{\chi}{\sqrt{1-\chi^2}} \quad (3)$$

$$\Delta_Z \equiv \frac{M_{\hat{Z}'}^2}{M_{Z_0}^2} \quad (4)$$

$$M_{Z_0}^2 \equiv \frac{1}{4}(g^2 + g'^2)v^2 \quad (5)$$

the mass eigenvalues of the Z and Z' (the photon remains massless) are [14],

$$M_Z^2 = \frac{M_{Z_0}^2}{2} \left[1 + s_W^2 \eta^2 + \Delta_Z \pm \sqrt{(1 - s_W^2 \eta^2 - \Delta_Z)^2 + 4s_W^2 \eta^2} \right] \quad (6)$$

$$M_{Z'}^2 = \frac{M_{Z_0}^2}{2} \left[1 + s_W^2 \eta^2 + \Delta_Z \mp \sqrt{(1 - s_W^2 \eta^2 - \Delta_Z)^2 + 4s_W^2 \eta^2} \right] \quad (7)$$

$$(8)$$

where the upper (lower) signs are for $\Delta_Z < 1 - s_W^2 \eta^2$ ($\Delta_Z > 1 - s_W^2 \eta^2$), with s_W the sine of the weak mixing angle. For consistency with precision data, we require $\eta \ll 1$ for which $M_Z \approx M_{Z_0}$ and $M_{Z'} \approx M_{\hat{Z}'}$.

The relevant couplings are given by:

$$g_\nu^Z = (\hat{g}_R^\nu P_R + \hat{g}_L^\nu P_L) s_\alpha \sqrt{1 + \eta^2} \quad (9)$$

$$g_\nu^{Z'} = (\hat{g}_R^\nu P_R + \hat{g}_L^\nu P_L) c_\alpha \sqrt{1 + \eta^2} \quad (10)$$

$$g_{\psi_{SM}}^{Z'} = -\frac{g}{c_W} c_\alpha (t_\alpha + \eta s_W) \left[T_{3L} P_L - s_W^2 Q \frac{t_\alpha + \eta/s_W}{t_\alpha + s_W \eta} \right] \quad (11)$$

$$g_{\psi_{SM}}^Z = -\frac{g}{c_W} c_\alpha (1 - t_\alpha \eta s_W) \left[T_{3L} P_L - s_W^2 Q \frac{1 - t_\alpha \eta/s_W}{1 - t_\alpha s_W \eta} \right] \quad (12)$$

with,

$$t_\alpha = \frac{-2\eta s_W}{1 - s_W^2 \eta^2 - \Delta_Z \pm \sqrt{(1 - s_W^2 \eta^2 - \Delta_Z)^2 + 4\eta^2 s_W^2}}. \quad (13)$$

¹ χ can be engineered to vanish in the UV, for example, by embedding $U(1)'$ into a larger gauge group which breaks down at scales of order $M_{\hat{Z}'}$.

The parameters defining the model thus include the WIMP and \hat{Z}' masses M_ν and $M_{\hat{Z}'} \approx M_{Z'}$, the mixing parameter $\eta \approx \chi$, and the Z' couplings to ν and to right-handed top quarks, $g_\nu^{Z'} \approx \hat{g}_\nu^{Z'}$ and $g_t^{Z'} \approx \hat{g}_t^{Z'}$. In addition we inherit the Higgs mass m_h as an unknown parameter of the Standard Model itself. Note that all the phenomenology of the model is controlled by the couplings of Z' to dark matter and to the top, and not through a direct coupling of dark matter to the Higgs [15].

As written, the model contains all of the most important features to describe dark matter, but is not UV complete. It contains $U(1)'^3$ and mixed $U(1)'$ -SM gauge anomalies which need to be cancelled, and the top Yukawa coupling is not $U(1)'$ invariant. One can imagine a variety of possible UV completions, usually from additional massive fermions, whose presence is not expected to significantly affect the phenomenology of interest here (see Appendix A for one particular example). Such fermions will also contribute to χ , and motivate our treatment of it as an independent parameter. We have also neglected the Higgs sector which breaks the $U(1)'$ symmetry, giving mass to the \hat{Z}' , as it is largely unimportant for the phenomenology of interest here.

As a result of its incompleteness, our theory is not renormalizable, and it is important to understand which quantities can be reliably computed in perturbation theory, and which are sensitive to the unspecified UV physics. At one-loop, four interactions are of particular importance to describe the production of gamma rays from WIMP annihilation, the effective Z' - γ - h , Z' - γ - Z , and Z' - γ - Z' vertices, all mediated through loops of top quarks, and the effective Z' - b - \bar{b} vertex, mediated by a loop of top quarks and W bosons. Naive power counting (confirmed by our explicit computations, see Section 3.2) combined with the demands of Lorentz and $SU(2) \times U(1)$ gauge invariance indicates that the Z' - γ - h , Z' - γ - Z , and Z' - γ - Z' vertices are all finite, whereas the Z' - b - \bar{b} vertex is log divergent, and thus sensitive to the details of the UV completion. While its precise value is thus ill-defined in our effective theory framework, we expect it to be of order,

$$g_{Z'b\bar{b}} \sim g_t^{Z'} \frac{\alpha}{32\pi s_W^2} \frac{m_t^2}{M_W^2} \log\left(\frac{\Lambda^2}{Q^2}\right) \quad (14)$$

where α is the fine-structure constant, M_W is the W boson mass, Λ is the scale of the UV completion of the theory², and Q^2 is a typical momentum transfer in the process of interest. Where relevant, we will use this estimate (with $\log \Lambda^2/Q^2 = 1$) as the actual value of the Z' - b - \bar{b} coupling, below.

2.2 Connection to RS Models

Our setup has a natural connection to RS theories in which the SM lives in the bulk [16], and such theories provide a natural UV completion. In RS theories, the hierarchy between the Planck and electroweak scales is explained through warping of an extra dimension, with the Higgs living on the IR boundary where the natural scale of physics is \sim TeV. As a result

²Given the need to generate the large top Yukawa interaction from higher dimensional operators at scale Λ , this scale cannot be much larger than $M_{Z'}$ itself without the effective theory breaking down. The log in Equation (14) is thus at most of order a few.

of the localized Higgs, the zero mode of the right-handed top quark must also live close to the IR brane, in order to realize the large top mass³.

As a consequence of the warping, all of the low level KK modes have wave functions whose support is concentrated near the IR brane. Thus, they inevitably couple strongly to t_R and the Higgs. If the dark matter particle is also a KK mode, for example a field whose boundary conditions forbid the appearance of a zero mass state, it will also couple strongly to the other KK modes. Taken all together, these features establish the main features of our effective theory. To connect more precisely to specific models, one should identify the DM and Z' fields as KK modes of bulk fields with the right properties.

Through the RS/CFT correspondence [17,18], the extra-dimensional theory is thought to be dual to an approximately scale-invariant theory in which most of the Standard Model is fundamental, but with the WIMP, Higgs, and right-handed top largely composite. The Higgs couples strongly to composite fields, and the amount of admixture in a given SM fermion determines its mass [19]. In this picture, the Z' is one of the higher resonances, built out of the same preons as the WIMP and t_R . RS theories provide a very motivated picture of the UV physics, but more generically, in any theory (not necessarily with approximate scale invariance) containing composite WIMPs [20] and composite top quarks [21] belonging to a common sector [22], one would expect strong couplings between them as a residual of the strong force which binds them, and perhaps negligible coupling to the rest of the Standard Model.

Some RS constructions automatically contain SM gauge singlet bulk fermions which can be identified with ν . The most obvious exists in models with an $SO(10)$ GUT symmetry in the bulk, for which the SM matter fits into 16 representations, including a gauge singlet with $(-, +)$ boundary conditions whose mass is generally somewhat atypically low compared to the other KK modes [13]. In this model, dark matter is stable as a result of a global symmetry needed to protect against too-rapid proton decay. Other possibilities include cases in which the dark matter is a stable neutral component [24,23] of a doublet under $SU(2)_L \times SU(2)_R$ (introduced to control precision corrections to ΔT [16]), though such objects have full strength $SU(2) \times U(1)$ interactions which we have not considered in this work.

In the $SO(10)$ model, the Z' represents the lowest KK mode of the $U(1)$ contained in $SO(10)$. It typically has mixing with the electroweak bosons, resulting in strong constraints from precision data. This will also be the case when the Z' is a KK mode of the electroweak bosons. As we will see, heavy Z' s still lead to interesting signals, and can satisfy relic density and direct detection constraints, provided their couplings are strong enough. We circumvent these constraints by considering a Z' whose mixing with the Z is kinetic. At large Z' masses this is not operationally different from the mass-mixing case, but it allows us to consider lower mass Z' s which are not ruled out by precision data.

As a final comment, the $SO(10)$ model also contains an electrically charged color triplet vector boson X_s^μ which couples directly to top and ν , and has mass of order the typical KK scale. We have chosen not to include X_s^μ in our effective theory, because it never approaches resonant behavior in WIMP annihilations, and thus is usually subleading compared to the \hat{Z}' .

³The left-handed top is usually chosen to be further from the IR brane, in order to mitigate constraints from precision electroweak tests [16].

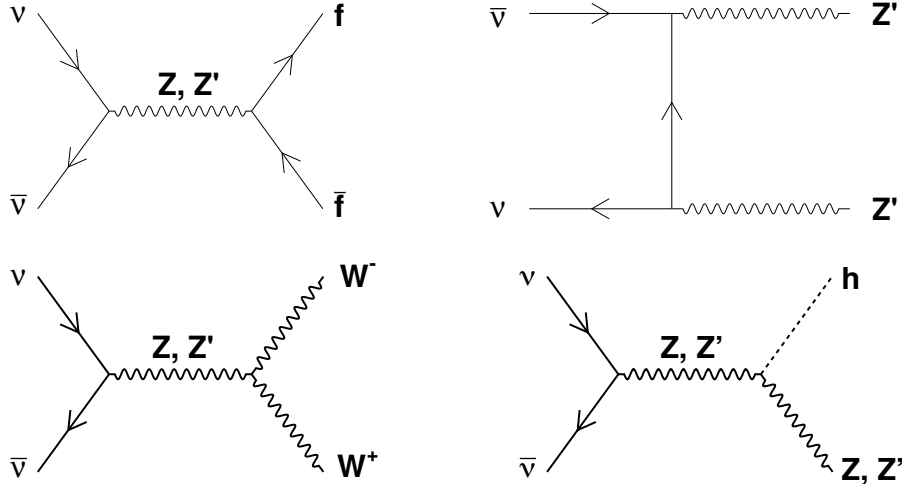


Figure 1: Representative tree-level Feynman diagrams for $\nu\bar{\nu}$ annihilation channels. If kinematically allowed, the two dominant channels are the s -channel Z' annihilation into $t\bar{t}$ and the t -channel annihilation into Z' . All other processes are suppressed by the kinetic mixing η or are loop-level.

3 Gamma Ray Spectrum from DM Annihilation

3.1 Continuum Photon Emission

The continuum photon emission largely originates from the decays of π^0 produced by hadronization of strongly interacting states and from gamma-rays emitted by light charged particles. As alluded to in the introduction, the resulting spectrum contains only very subtle information about the primary products of dark matter annihilation, as their shapes are very similar. In the particular case of $\nu\bar{\nu}$ annihilation, the dominant processes are those shown in Fig. 1.

One interesting aspect of our model is that it leads to suppressed continuum emission, which increases the prominence of the gamma ray lines. For M_ν less than both m_t and $M_{Z'}$, the natural annihilation into top pairs is closed, forcing annihilation predominantly into light SM particles, whose rates are suppressed by the small kinetic mixing η or the loop-induced Z' - $b\bar{b}$ interaction. For M_ν greater than either m_t or $M_{Z'}$, $\nu\bar{\nu}$ annihilation is dominantly into top quark pairs, and the continuum emission is more sizable. Nevertheless, since the continuum generated by $t\bar{t}$ annihilation is softer for $M_{DM} \lesssim 200$ GeV than the one obtained from $b\bar{b}$ or WW annihilation, even in this case, we expect reduced continuum emission close to the cut-off of the spectrum, where the lines are located. In Fig. 2 we show the spectra for generic annihilation into $b\bar{b}$, WW or $t\bar{t}$ and spectra for $\nu\bar{\nu}$ annihilation for two parameter sets with masses above and below m_t , both of which lead to the correct thermal relic density (see Section 4.2). The photon fluxes induced by $\nu\bar{\nu}$ annihilations have been computed for a Navarro-Frenk-White DM density profile and refer to an observation of the galactic central region with an angular acceptance $\Delta\Omega = 10^{-5}$ sr (see Section 3.4 for more details).

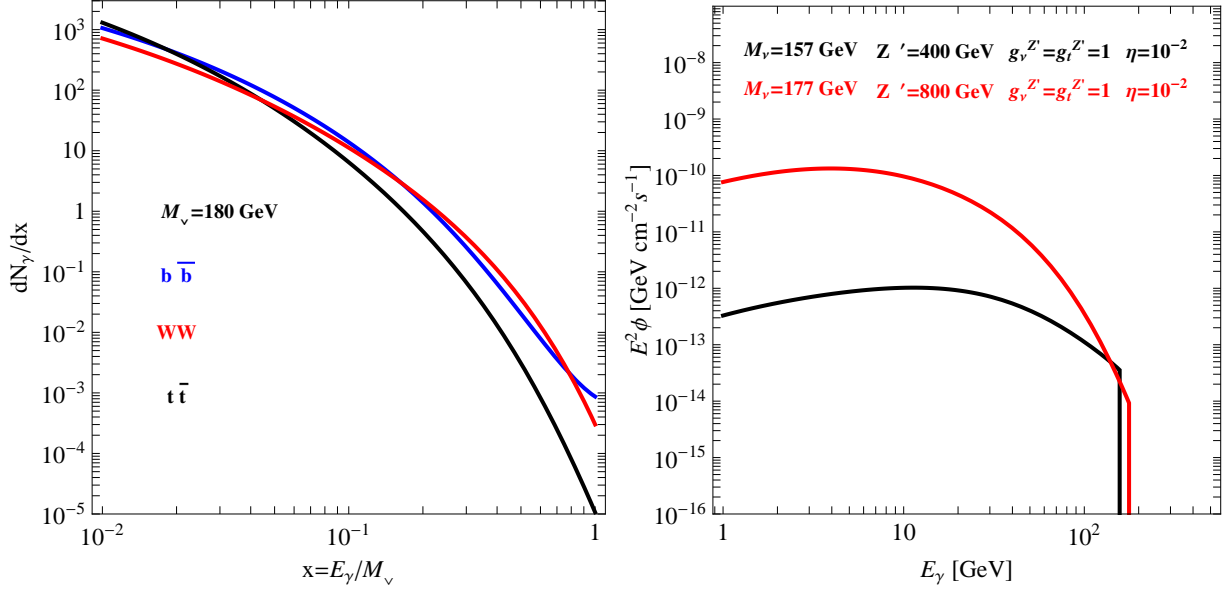


Figure 2: (Left panel) Continuum photon spectrum dN_γ/dx , with $x = E_\gamma/M_{DM}$, obtained from annihilation into $b\bar{b}$, WW or $t\bar{t}$. (Right panel) Comparison of the continuum spectra obtained for two parameter sets with $M_\nu < m_t$ and $M_\nu > m_t$. Fluxes are for an observation of a $\Delta\Omega = 10^{-5}$ sr region around the galactic center. A NFW density profile has been employed.

3.2 Gamma Ray Lines

In this section we compute the expected gamma ray line intensities, for the γh , γZ and $\gamma Z'$ lines. An interesting feature that results from an s -channel Z' being the sole portal from the WIMP sector to the SM is the fact that there is no $\gamma\gamma$ line, as dictated by the Landau-Yang theorem [25] (see also [26] for a more recent discussion). At leading order in the WIMP relative velocity, the cross section into γX is given by,

$$\sigma v = \frac{1}{64\pi M_\nu^2} \left(1 - \frac{M_X^2}{4M_\nu^2}\right) \overline{|\mathcal{M}|^2}. \quad (15)$$

In the sections below, we outline the computation of the the matrix elements \mathcal{M} for annihilation into γh , γZ and $\gamma Z'$.

3.2.1 γh

A heavy neutrino and anti-neutrino can annihilate into a γh final state through a Z' boson which connects to a loop of top quarks (see Fig. 3). In the non-relativistic limit, the matrix element, averaged over initial heavy neutrino spins and summed over the final state photon polarizations may be expressed as,

$$\overline{|\mathcal{M}|^2}_{\gamma h} = \frac{\alpha\alpha_t N_c^2}{72\pi^2} \mathcal{V}^2 \frac{(g_t^{Z'} g_\nu^{Z'})^2 M_\nu^2 m_t^2}{(4M_\nu^2 - M_{Z'}^2)^2 + M_{Z'}^2 \Gamma_{Z'}^2} \quad (16)$$

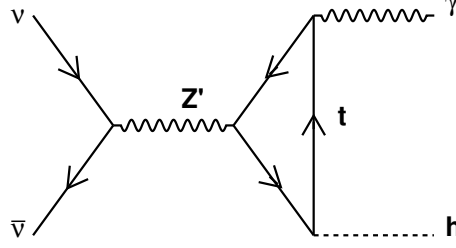


Figure 3: Representative Feynman diagram for $\nu\bar{\nu} \rightarrow \gamma h$.

where $N_c = 3$ is the number of colors, $\alpha_t = y_t^2/(4\pi)$ is the fine structure constant corresponding to the top Yukawa interaction, and $\Gamma_{Z'}$ is the Z' width.

The vertex factor \mathcal{V} describes the effective coupling between an off-shell Z' boson and on-shell Higgs and photon, and may be expressed as,

$$\begin{aligned} \mathcal{V}_{\gamma h} &= \frac{8M_\nu^2}{(m_h^2 - 4M_\nu^2)} [B_0(m_h^2; m_t, m_t) - B_0(4M_\nu^2; m_t, m_t)] \\ &+ [4M_\nu^2 + 4m_t^2 - m_h^2] C_0(m_h^2, 0, 4M_\nu^2; m_t, m_t, m_t) + 2 \end{aligned} \quad (17)$$

where the scalar integrals are defined as,

$$\begin{aligned} B_0(p^2; m, m) &= 16\pi^2 \int \frac{d^n \ell}{(2\pi)^n} \frac{1}{\ell^2 - m^2} \frac{1}{(\ell + p)^2 - m^2}, \quad (18) \\ C_0(p_a^2, p_b^2, (p_a + p_b)^2; m_1, m_2, m_3) &= 16\pi^2 \int \frac{d^n \ell}{(2\pi)^n} \frac{1}{\ell^2 - m_1^2} \frac{1}{(\ell + p_a)^2 - m_2^2} \frac{1}{(\ell + p_a + p_b)^2 - m_3^2}. \end{aligned}$$

3.2.2 γZ and $\gamma Z'$

The calculation of the γZ and $\gamma Z'$ cross sections follows along the same lines as that of γh , but is somewhat more complicated by the Z (or Z') spin indices. In this subsection, we outline the calculation of the cross sections for $\nu\bar{\nu} \rightarrow \gamma Z(Z')$ and reserve detailed expressions for the appendix.

We begin with the effective vertex depicted in Fig. 4. The expression for the effective vertex is given by:

$$\Gamma^\alpha = \epsilon_\mu^*(p_A) \epsilon_\nu^*(p_Z) \sum_{i=1}^5 C_i M_i^{\alpha\mu\nu}. \quad (19)$$

where p_A and p_Z are the momenta of the photon and the outgoing Z (Z'), respectively, and we have accounted for the transversality of the polarization tensors ($\epsilon^*(p_A) \cdot p_A = \epsilon^*(p_Z) \cdot p_Z = 0$). The C_i coefficients are functions of the top quark couplings to Z and Z' as well as the scalar integrals given in Eqs. (18) with the replacement $m_h \rightarrow M_Z$. The exact expressions for these coefficients are given in appendix B. The $M_i^{\alpha\mu\nu}$ tensor structures are expressed in terms of

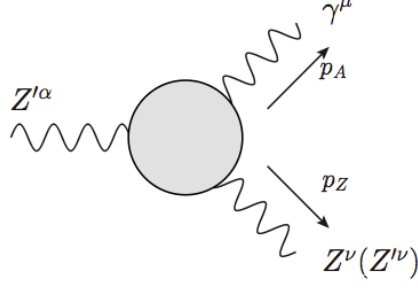


Figure 4: Labels for the Z' - γ - Z/Z' effective vertex.

Levi-Civita tensors as:

$$M_1^{\alpha\mu\nu} = p_{A,\lambda}\epsilon^{\lambda\mu\nu\alpha}, \quad (20)$$

$$M_2^{\alpha\mu\nu} = p_{Z,\lambda}\epsilon^{\lambda\mu\nu\alpha}, \quad (21)$$

$$M_3^{\alpha\mu\nu} = p_A^\alpha p_{A,\lambda} p_{Z,\sigma} \epsilon^{\lambda\sigma\mu\nu}, \quad (22)$$

$$M_4^{\alpha\mu\nu} = p_Z^\alpha p_{A,\lambda} p_{Z,\sigma} \epsilon^{\lambda\sigma\mu\nu}, \quad (23)$$

$$M_5^{\alpha\mu\nu} = p_A^\nu p_{A,\lambda} p_{Z,\sigma} \epsilon^{\lambda\sigma\alpha\mu}. \quad (24)$$

After summing over the polarizations of the external gauge bosons, only M_1 , M_2 and M_5 actually contribute to the matrix element squared.

Coupling the effective vertex Γ^α to the $Z'\nu\bar{\nu}$ vertex and averaging (summing) over initial (final) state spins and polarizations, we find the matrix-element-squared for $\nu\bar{\nu} \rightarrow \gamma Z$ is given by:

$$|\overline{\mathcal{M}}|^2_{\gamma Z} = \frac{\alpha^2 N_c^2}{576\pi^2 s_W^2 c_W^2} \mathcal{V}_{\gamma Z}^2 \frac{(g_\nu^{Z'})^2 M_\nu^2}{(4M_\nu^2 - M_{Z'}^2)^2 + M_{Z'}^2 \Gamma_{Z'}^2}, \quad (25)$$

where $s_W(c_W)$ is the sine (cosine) of the SM weak-mixing angle. The matrix-element-squared for $\nu\bar{\nu} \rightarrow \gamma Z'$ is:

$$|\overline{\mathcal{M}}|^2_{\gamma Z'} = \frac{\alpha N_c^2}{144\pi^3} \mathcal{V}_{\gamma Z'}^2 \frac{(g_\nu^{Z'})^2 M_\nu^2}{(4M_\nu^2 - M_{Z'}^2)^2 + M_{Z'}^2 \Gamma_{Z'}^2}. \quad (26)$$

The expressions for $\mathcal{V}_{\gamma Z}^2$ and $\mathcal{V}_{\gamma Z'}^2$ are given in the appendix. The respective cross sections are then obtained by substituting Eqs. (25) and (26) into Eq. (15). In Fig. 5, we plot the cross sections for γh , γZ and $\gamma Z'$ as a function of the neutrino mass M_ν for $g_\nu = g_t = 3$ and for two values of $M_{Z'}$. We compare them with the continuum obtained from all 2-body annihilations.

3.3 Gamma Ray Spectra

The photon spectrum originating from the process $\nu\bar{\nu} \rightarrow \gamma X$ deviates from a monochromatic emission due to the finite decay width Γ_X of the unstable particle X with mass M_X and is given by [3]:

$$\frac{dN_\gamma^X}{dE} = \frac{4M_\nu M_X \Gamma_X}{f_1 f_2} \quad (27)$$

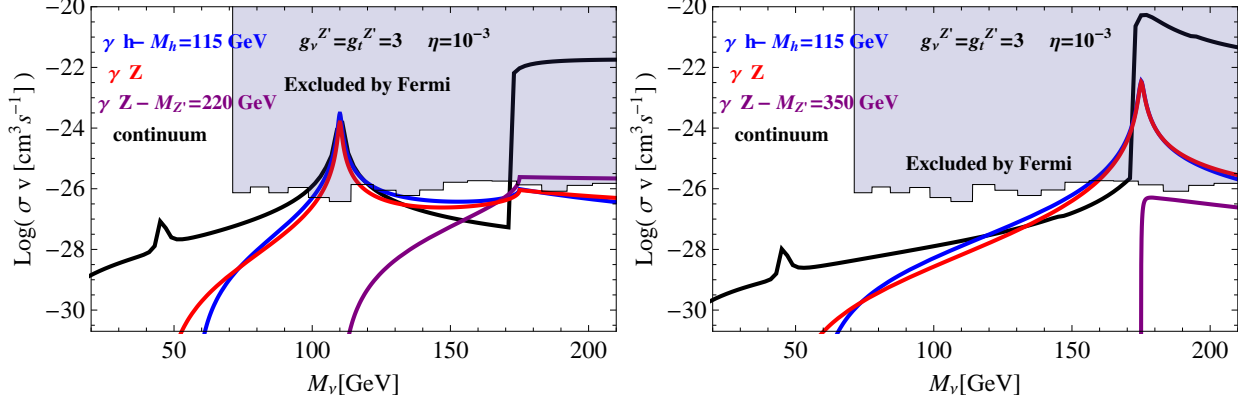


Figure 5: Cross sections as a function of the WIMP mass for two values of $M_{Z'}$ and some choice of couplings for the three gamma line emission channels γh (blue curves), γZ (red curves) and $\gamma Z'$ (purple curves) along with the continuum γ ray emission (black curves). Also shown are 95% C.L. exclusion limits inferred from Fermi data [27] for a γh line with $M_h = 115$ GeV and a NFW profile.

where

$$\begin{aligned}
 f_1 &= \tan^{-1} \left(\frac{M_X}{M_\nu} \right) + \tan^{-1} \left(\frac{4M_\nu^2 - M_X^2}{M_X \Gamma_X} \right) \\
 f_2 &= (4M_\nu^2 - 4M_\nu E_\gamma - M_X^2)^2 + \Gamma_X^2 M_X^2.
 \end{aligned}$$

The Higgs width Γ_h will match the SM prediction at tree level, but at loop level could have tiny contributions from $h \rightarrow \gamma Z'$ or $h \rightarrow Z' Z'$, depending on m_h and $M_{Z'}$. We neglect these insignificant corrections in the inclusive width. The tree level Z' width is shown in Fig. 6, although the final results are not sensitive to its precise value except very close to the Z' resonance.

The detection of gamma ray lines *per se* represents smoking-gun evidence for dark matter annihilation, but it does not tell us which processes are responsible for the observed lines. However, additional indirect dark matter searches, direct detection experiments, and LHC observations can complement the information from gamma-ray telescopes. For example, the energies of gamma-lines probe the masses of the particles in the associated annihilation process, c.f. Eq. (1), and this could be combined with independent measurements of particle masses at colliders. This cross-check could prove extremely useful to identify a given long-lived particle produced at colliders as a significant fraction of the dark matter present in the galaxy.

The detection and identification of the $Z\gamma$ and $h\gamma$ lines could also allow one to determine the Higgs mass. Fig. 7 shows the region in the M_ν - m_h plane where these two lines are potentially separately observable. The $h\gamma$ line can be distinguished from the $Z\gamma$ line if the energy separation is at least twice the energy resolution of the experiment, which for the Fermi LAT is $\delta E \sim 10\%$ for the energies of interest [28]. The maximum Higgs mass which can be probed in $\nu\bar{\nu} \rightarrow h\gamma$ annihilation is $2M_\nu$. For $2M_\nu > M_{Z'}$, the $Z'\gamma$ line is also present. In Fig. 7 we show, for the representative cases of $M_{Z'} = 220$ GeV and 500 GeV, the combination of Higgs and ν masses for which all three of the lines are distinguishable by an experiment with $\sim 10\%$ energy resolution.

$$g_v^{Z'} = g_t^{Z'} = 1$$

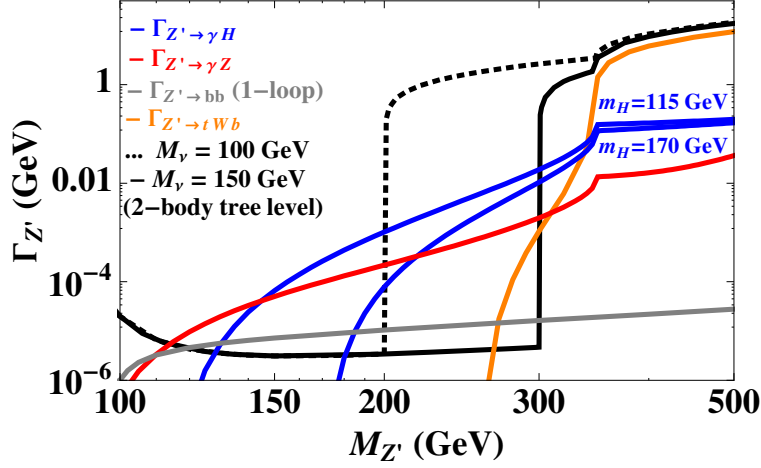


Figure 6: Z' partial widths as a function of its mass for two values of the WIMP mass. The solid and dotted black lines are for the sum of the partial widths into all 2-body tree level modes, which, apart from the $t\bar{t}$ channel, are induced by the kinetic mixing η . The gray line is for the one-loop annihilation into $b\bar{b}$, estimated using Eq. 14. One-loop annihilation into γh and γZ are shown in blue and red respectively. The orange line is for the three-body annihilation into tWb .

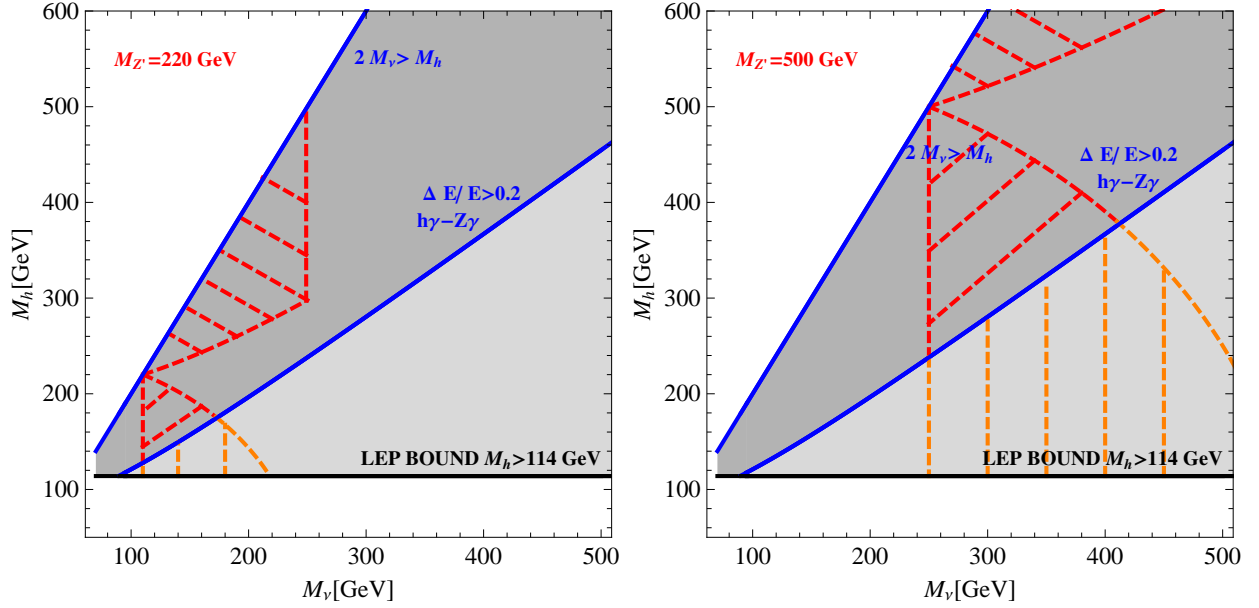


Figure 7: Regions of the M_ν - m_h parameter space (for $M_{Z'} = 220$ GeV/500 GeV in the left/right panels) for which the γZ and γh lines can be distinguished by an experiment with 10% energy resolution (dark grey); in the light grey region they are merged. The red dashed area further shows the regions where the γZ , γh , and $\gamma Z'$ lines are distinguishable. In the dashed orange region γZ and γh lines are merged but distinguishable from the $\gamma Z'$ line.

In most DM models producing line signals, there will typically be one line from annihilation into $\gamma\gamma$ and/or γZ (these two lines can be resolved only if the wimp mass is in the $\sim 50 - 100$ GeV mass range). Measuring the energy of this line will provide useful information on the DM mass. In addition, if another less energetic line is detected, this will allow to estimate the mass of the new heavy particle X that DM annihilates into. Since we are considering gamma ray energies between a few GeV to a few hundreds of GeV, this means that the DM and X particles will be kinematically accessible at the LHC (if heavier than a TeV the corresponding gamma ray signal will be suppressed). Therefore any line signal observed with FERMI, MAGIC or HESS should be accompanied by a signal at the LHC.

3.4 Dark Matter Distribution

The differential photon flux produced by dark matter annihilations and collected from a region of angular size $\Delta\Omega$ is computed as:

$$\Phi_\gamma(E_\gamma) = \frac{1}{4\pi} \frac{r_\odot \rho_\odot^2}{4M_\nu^2} \frac{dN_\gamma}{dE_\gamma} \bar{J} \Delta\Omega \quad (28)$$

with

$$\begin{aligned} \frac{dN_\gamma}{dE} &= \sum_f \langle \sigma v \rangle_f \frac{dN_\gamma^f}{dE} & \bar{J} &= \frac{1}{\Delta\Omega} \int_{\Delta\Omega} J(\psi) \\ J(\psi) &= \int_{los} \frac{ds}{r_\odot} \left(\frac{\rho(r(s, \psi))}{\rho_\odot} \right)^2. \end{aligned} \quad (29)$$

dN_γ/dE includes all possible annihilation final states f , with $\langle \sigma v \rangle_f$ and dN_γ^f/dE referring to the corresponding cross sections and photon spectra per annihilation. The factor $1/4$ in Eq. (28) is appropriate for a Dirac fermion WIMP with predominantly particle-anti-particle annihilation modes. The dimensionless quantity $J(\psi)$ corresponds to the integration of the photon signal along a line of sight making an angle ψ with the direction of the galactic center. The total observed flux is then obtained integrating the emission over the the observed region of angular size $\Delta\Omega$. The normalization factors $\rho_\odot = 0.3 \text{ GeV cm}^{-3}$ and $r_\odot = 8.5 \text{ kpc}$ correspond respectively to the dark matter density at the solar position and to the distance of the Sun from the galactic center. We model the dark matter density distribution in our galaxy, $\rho(x)$, as a Navarro-Frenk-White (NFW) profile [29], which is a good fit to current N -body simulations:

$$\rho_{\text{NFW}}(r) = \frac{\rho_s}{\frac{r}{r_s} \left(1 + \frac{r}{r_s} \right)^2}. \quad (30)$$

Some recent simulations, however, prefer the so-called ‘‘Einasto’’ [30] profile, which is slightly more shallow at small radii and it does not converge to a definite power-law:

$$\rho_{\text{Einasto}}(r) = \rho_s \cdot \exp \left[-\frac{2}{\alpha} \left(\left(\frac{r}{r_s} \right)^\alpha - 1 \right) \right], \quad \alpha = 0.17. \quad (31)$$

MW halo model	r_s in kpc	ρ_s in GeV/cm ³	\bar{J} (10^{-5})
NFW [29]	20	0.26	$15 \cdot 10^3$
Einasto [30]	20	0.06	$7.6 \cdot 10^3$
Adiabatic [31]			$(1.2 - 14.0) \cdot 10^7$

Table 1: Parameters of the dark matter density profiles for the Milky Way discussed in the text and corresponding value of \bar{J} for $\Delta\Omega = 10^{-5}$.

The presence of baryons, not accounted for in the simulations previously quoted, may significantly change the picture, particularly in the inner region of the galaxy where the gravitational influence of the super massive black hole is expected to have a large feedback on the surrounding dark matter distribution. The evolution of the dark matter density profile, accounting for dark matter-star interactions, capture in the central black-hole and the presence of dark matter annihilations, has been simulated in the so-called ‘‘adiabatic compression’’ scenario in Ref. [31]. The final density distribution is significantly increased at small radii with respect to the initial NFW profile.

In Table 1 we show the \bar{J} factor for the dark matter distributions discussed here for an observation of the galactic central region with an angular acceptance $\Delta\Omega = 10^{-5}$, corresponding to the angular resolution of Fermi-LAT and current Air Cherenkov Telescopes (ACTs). The adiabatically contracted profiles of Ref. [31] depend on the wimp annihilation cross section and in Table 1 we list the range of \bar{J} factors that have been used in Fig.8. The large uncertainties in the dark matter distribution in the region considered turn into large uncertainties on the predicted photon fluxes. On the other hand the ρ^2 dependence of the signal suggests the galactic central region as the best target to maximize the signal. For the rest of the paper we adopt NFW as a dark matter profile benchmark and for illustration we show some of the predictions for the adiabatically-contracted profile.

3.5 Predicted Photon Fluxes and Comparison with Experiments

The expected photon signal, Φ_γ^S is obtained by convolving the photon flux in Eq. 28 with the energy response of the instrument $G(E_0, E)$:

$$\Phi_\gamma^S(E) = \int dE_0 \Phi_\gamma(E_0) G(E_0, E) \quad (32)$$

where we assume a Gaussian kernel

$$G(E_0, E) = \frac{1}{\sqrt{2\pi}E\sigma} \exp\left(-\frac{(E_0 - E)^2}{2\sigma^2 E^2}\right) \quad (33)$$

with σ depending on the detector energy resolution ξ as $\sigma = \xi/2.3$.

In Fig. 8, we show the predicted photon fluxes at the galactic center for different choices of particle physics parameters which give the correct thermal relic abundance and satisfy the constraints from direct detection. For comparison we plot the HESS observations of the same angular region [32] and the EGRET data on the unidentified source 3EG J1746-2851 [33,34], corresponding instead to $\Delta\Omega = 10^{-3}$, appropriate for a detector with an angular resolution of

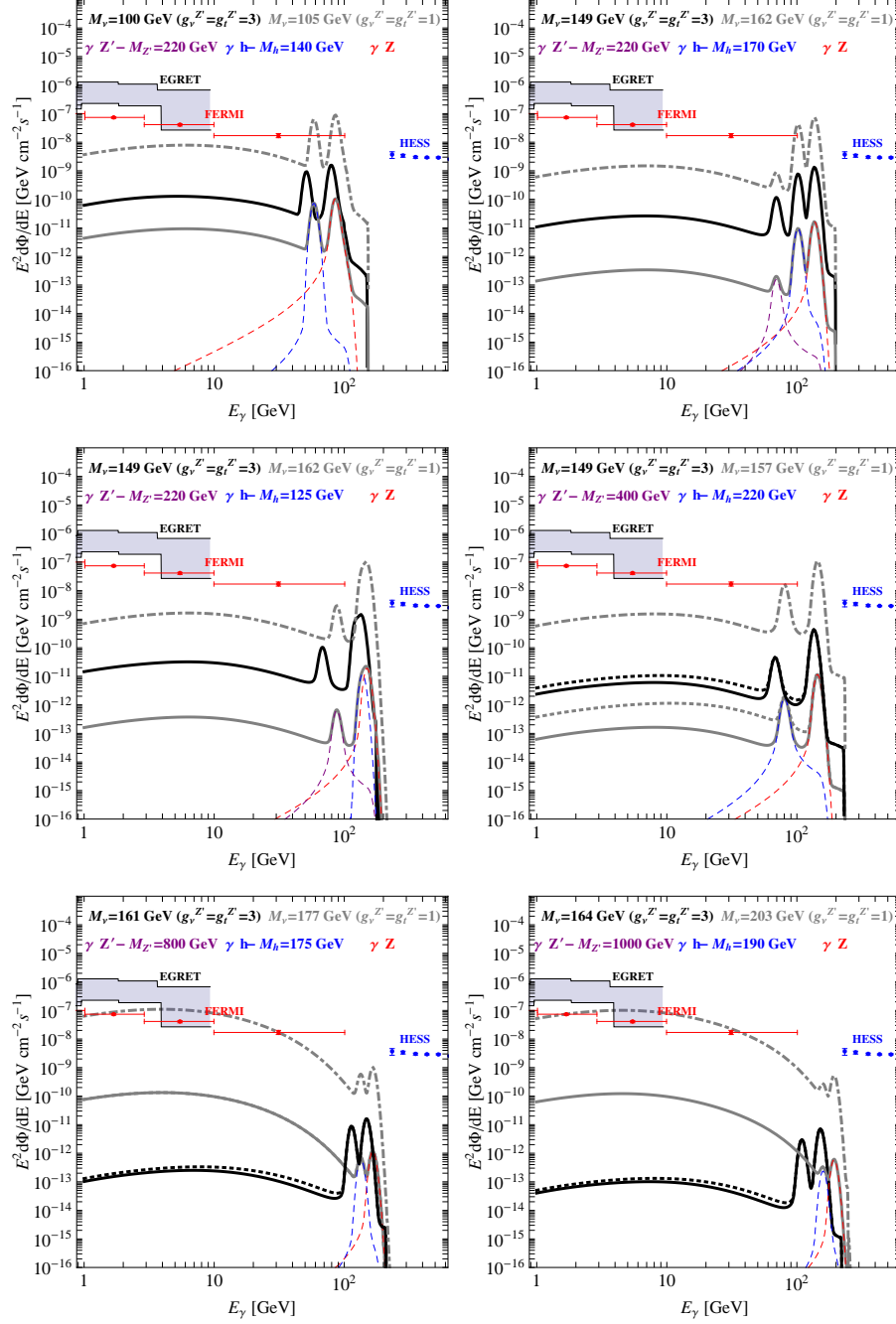


Figure 8: Spectra obtained for different choices of mass parameters and coupling $g_{\nu,t}^{Z'} = 1$ (black), $g_{\nu,t}^{Z'} = 3$ (gray), $\eta = 10^{-3}$ (solid), $\eta = 10^{-2}$ (dotted) that lead to the correct relic density and satisfy direct detection constraints. Upper plots are for $\eta = 10^{-3}$ only since for these choices of couplings and $M_{Z'}$ mass $\eta = 10^{-2}$ is excluded by direct detection constraints. $\Delta\Omega = 10^{-5}$, and a NFW dark matter profile is assumed. Dot-dashed lines are for the adiabatically-contracted profile in Table 1, $\Delta\Omega = 10^{-5}$, $g_{\nu,t}^{Z'} = 1$ and $\eta = 10^{-3}$. EGRET data are from [33, 34], HESS from [32] and Fermi from [35].

$\sim 1^\circ$. Fermi satellite preliminary results fill the region in energy between HESS and EGRET, providing the most powerful probe of WIMP annihilation into gamma rays to date.

The Fermi collaboration has recently presented preliminary 95% C.L. upper limits on the line cross section from a line search based on a shape analysis [35] of 11 months of observation of a region excluding the galactic plane but including the galactic central region. We translate the preliminary bounds on $\gamma\gamma$ and γZ to bounds on γX for $M_X = 115$ GeV. Comparing even the current bounds with the γh cross sections in Fig. 5, we see that Fermi is *already* ruling out a small region of parameter space where the annihilation is close to on-resonance. Thus, even though its line search so far has had a null result, longer exposure and optimized windows of observation still have major potential to observe a signal. Emission from dark matter is expected to be maximized at the galactic center, but that may not be the optimal region to detect this signal. For example, the intense point-like gamma-ray source at the galactic center detected by the ACTs at energies above 200 GeV and confirmed by Fermi at lower energies, constitutes a serious background for dark matter searches. The morphology of the background and signal emissions suggests that a better strategy could be to consider a larger angular region, with the size depending on the choice of dark matter profile (e.g. [36]) and subtract all the astrophysical point sources there detected. This kind of analysis is ongoing within the Fermi collaboration [27].

In Fig.8 we are not exhausting all the possible predictions of our model. The spectra shown are instead meant to be illustrative cases which capture different phenomenological aspects of the theory and present some of the possible outcomes rather than characterize all of them. For instance, we see how, depending on the relative masses of the WIMP and the Higgs, there can be one, two (either $\gamma-h$ and $\gamma-Z$ or if these two are merged, $\gamma-Z'$ and $\gamma-h/Z$) or three lines and for WIMP masses below the top mass, as it happens in the large coupling regime, there are spectra with prominent lines compared to the continuum. Taking also into account the uncertainties in the dark matter distribution, which we illustrate by considering two different profiles, it is clear that the line signals we are presenting may be at the reach of future Fermi observations. Besides Fermi, the next generation of low-threshold, air cherenkov telescopes, such as CTA [37] may potentially have a dramatic impact on dark matter searches. In particular, they may probe the line signals shown in Fig.8 due to their improved energy threshold, which is expected to be lowered to ~ 30 GeV⁴.

4 Other Signals and Constraints

4.1 Elastic Scattering and Direct Detection

The fact that there has been no unequivocal observation to date of dark matter scattering with heavy nuclei places a constraint on any theory of dark matter and, in particular, results

⁴Considering a future CTA observation of a 2° region around the galactic center and based on a signal-to-noise analysis, we have sketched the 3σ evidence prospects (assuming a NFW profile) for a generic line γX . We have assumed a CTA effective area of ~ 1 km², an energy resolution of 15% and 200 hours of observations. The prospects strongly depend on the hadronic rejection efficiency, which we take energy independent and we vary from a conservative value of $\epsilon = 50\%$ to a more optimistic $\epsilon = 1\%$. For example, focusing on a γX line with $M_X = 115$ GeV and for $M_\nu = 100$ GeV the cross-section at the reach of CTA are $\sim 2 \cdot 10^{-28}$ cm³ s⁻¹ ($\epsilon = 1\%$) and $\sim 10^{-27}$ cm³ s⁻¹ ($\epsilon = 50\%$).

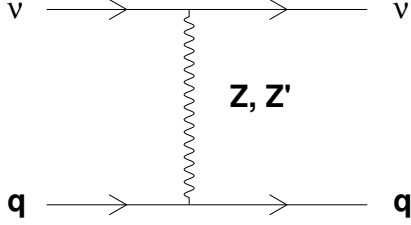


Figure 9: Representative diagram indicating how ν may scatter elastically with nuclei.

in a bound on the Z - Z' mixing parameter χ of Eq. (2). As shown in Fig. 9, ν scatters with nuclei largely by exchanging a Z or Z' boson which couples to valence quarks. This is a direct consequence of the fact that ν is a Dirac fermion, which (unlike a Majorana one) generally has vector interactions which remain large in the non-relativistic limit. As a result, the non-observation of a signal at direct detection experiments puts strong constraints on the couplings involved in the reaction. The ν -nucleon elastic scattering cross section is given by:

$$\sigma_{\substack{\text{neutron} \\ \text{proton}}} = \frac{g^2 m_{n/p}^2 g_\nu^{Z'}{}^2}{64\pi c_W^2} \left[\frac{t_\alpha c_\alpha}{M_Z^2} \left(1 - 4s_W^2 + 3\eta t_\alpha s_W \right) - \frac{A_\alpha}{M_{Z'}^2} \left(1 - 4s_W^2 B_\alpha \right) \right]^2, \quad (34)$$

where,

$$A_\alpha = c_\alpha(t_\alpha + \eta s_W) \quad \text{and} \quad B_\alpha = (t_\alpha + \eta/s_W)/(t_\alpha + \eta s_W). \quad (35)$$

Elastic scattering of ν with target nuclei is entirely on protons because the scattering on neutrons vanishes due to a cancellation between the Z and Z' contributions. The proton scattering is almost entirely from the Z' exchange.

In Fig. 10 we show the upper limit on the product of $g_\nu^{Z'} \times \eta$ based on the null result of the latest CDMS search for elastic WIMP-Germanium scattering [38]. The dependence on the ν mass is relatively mild provided it is greater than the mass of Germanium, about 70 GeV. Provided the Z' mass is larger than a few tens of GeV, the constraints are consistent with order one coupling between ν and the Z' , and η consistent with the expectation that it is induced at the one loop level. Note that these constraints are stronger than those from the EW precision measurements

$$\left(\frac{\eta}{0.1} \right)^2 \left(\frac{250 \text{ GeV}}{M_{Z'}} \right)^2 \lesssim 1, \quad (36)$$

derived in [14].

4.2 Relic Density

As a Dirac fermion, ν can carry a conserved $U(1)_\nu$ global charge, raising the possibility that its relic density may be understood as a primordial asymmetry between WIMP particles and antiparticles⁵. However, even in the absence of a $\nu - \bar{\nu}$ asymmetry, ν has roughly the right

⁵Of course, annihilation into gamma rays today requires a small breaking of $U(1)_\nu$ such that particles and antiparticles re-equilibrate after freeze-out.

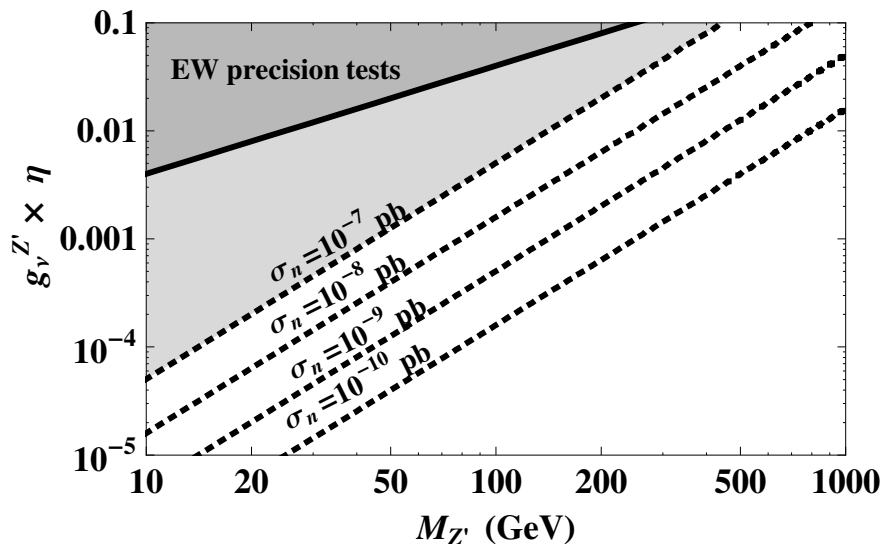


Figure 10: Contours for the ν -nucleon elastic scattering cross section. The dark grey region is excluded by EW precision data. The lighter grey region indicates the approximate region ruled out by CDMS (the precise limit depends on the WIMP mass).

properties for its thermal relic abundance to be appropriate to match cosmological measurements. Thus, while we present the relic density as a very interesting way to understand the parameter space, it is not a firm bound in the sense that the density of ν today may be due to a primordial asymmetry or a nonstandard cosmology. In fact, we do find regions with the correct thermal relic abundance which can lead to beautiful gamma ray line signatures at Fermi.

The abundance of a thermal relic ν is controlled by its annihilation cross section into SM particles, mediated by the same diagrams which control annihilation into gamma rays today shown in Fig. 1. To circumvent the constraints from elastic scattering in Fig. 10, we restrict our discussion to $M_{Z'} \gtrsim 200$ GeV, for which the results can be broadly classified:

- For $M_\nu \lesssim M_{Z'}/2$, the annihilation proceeds largely through the ν coupling to Z' , into light SM fermions and is controlled by the size of the kinetic mixing. The relic density is typically in conflict with direct detection constraints unless the annihilation is very close to the resonance, and the small WIMP mass kinematically forbids all of the line signals.
- For $M_{Z'}/2 \lesssim M_\nu \lesssim m_t$, annihilation is via loop-level processes into γZ , γh and $b\bar{b}$ (for $M_\nu \sim m_t$ there is also some annihilation into off-shell $t\bar{t}$). Far below the Z' resonance for modest ν coupling to the Z' , the rate is typically too small, leading to over-abundance of ν in the early Universe. This is remedied by having M_ν slightly above or below $M_{Z'}/2$ or large Z' couplings, for which the relic density prefers M_ν slightly smaller than m_t .
- For $M_\nu \gtrsim m_t$, the non-relativistic annihilation proceeds largely into a $t\bar{t}$ final state.

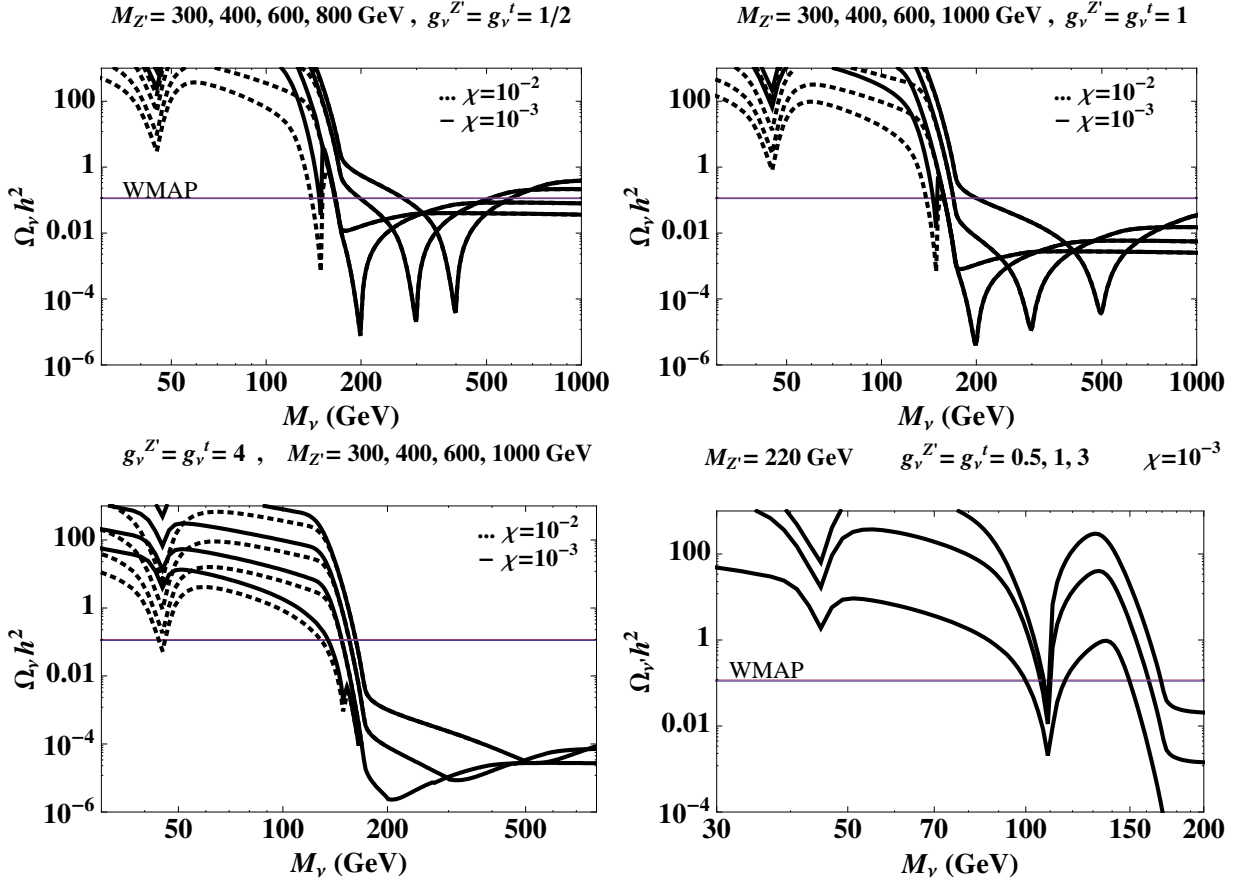


Figure 11: Relic density as a function of the ν mass for the indicated parameters. As the Z' coupling to top and ν increases, the prediction for the ν mass gets narrower assuming ν is the sole contributor to the relic density: For $g_{\nu}^{Z'}, g_t^{Z'} \gtrsim 1$ this implies $M \sim 150$ GeV, whatever the value of $M_{Z'} \gtrsim 400$ GeV. For lighter $M_{Z'}$, several M_{ν} can lead to the correct relic density. The calculation includes all 2-body tree level annihilations as well as the loop-induced annihilations into $b\bar{b}$.

For moderate couplings, $g_{\nu}^{Z'} \sim g_t^{Z'} \sim 1/2$, there is a continuum of possibilities which are highly correlated with the position of the Z' resonance. As the coupling is dialed stronger this window becomes narrower and occurs for ν masses close but smaller than the top mass. Taken together with the previous case, this is a robust prediction of the model in the strong coupling regime. Whatever the value of $M_{Z'} \gtrsim 350$ GeV, the DM mass is around 150 GeV, far away from the Z' resonance.

- Finally, if $M_{\nu} \gtrsim M_{Z'}$, the t -channel process $\nu\bar{\nu} \rightarrow Z'Z'$ opens up. However, this does not play an important role since we are considering $M_{Z'} \gtrsim 200$ GeV and in this case the annihilation cross section into top quark pairs continues to dominate.

The predicted thermal relic density for several representative parameter sets are shown in Fig. 11. The calculation was done with MicrOMEGAs [39].

4.3 Anti-matter Signals and radio constraints

Other than photons, ν annihilations produce e^+ and \bar{p} fluxes which can leave their imprint in the energy spectra of galactic cosmic-rays. PAMELA \bar{p} data can in particular set stringent constraints on such contributions, as it has been noticed in Refs. [40–42]. We have computed \bar{p} fluxes induced by tree level $\nu\bar{\nu}$ annihilations into SM particles as well as \bar{p} produced by h and Z decays associated with the γh and γZ line processes. For each model in Fig. 8 the antiproton signal is well below the experimental data. Generally, since relic density calculations usually predict ν masses of the order $\mathcal{O}(100)$ GeV, we don't expect any features in the antiproton spectrum at energies above present PAMELA measurements. The same considerations hold for cosmic-ray positrons: the predictions fall in the energy range explored by PAMELA and we find a very small ν contribution to the measured positron fraction, i.e. $e^+/(e^+ + e^-)$. While ν may not describe the PAMELA excess, several astrophysical processes such as pulsars [43] or supernovae remnants [44] offer an explanation. For illustration, in Fig. 12 we show the predictions for e^+ and \bar{p} signals for one model in Fig. 8. Inverse Compton scattering of high energy electrons and positrons (produced by ν annihilations) off interstellar photons produce gamma-rays. The constraints on such emission from EGRET data and preliminary Fermi results in Ref. [45,46] are easily satisfied by the models in Fig. 12, because of the small e^\pm fluxes associated with ν annihilations. Further bounds on the WIMP annihilation cross section are obtained comparing radio observations of the galactic center with the synchrotron emission induced by the propagation of such high energy electrons and positrons in the galactic magnetic field. The constraints are powerful but strongly dependent on the choice of the DM density profile, as shown in Refs. [42,47–49]. In particular, current N-body simulations cannot probe the DM density distribution in the inner region of the galaxy $< \mathcal{O}(100)$ pc). For a simple extrapolation of the NFW density profile suggested by simulations at these small radii, we checked that the combination of masses and cross-sections considered in Fig. 8 are compatible with the radio bounds. For the enhanced profile of Ref. [31], resulting from the DM adiabatic contraction around the central black hole of our galaxy, the gamma-ray fluxes can be in tension with the data and the associated synchrotron emissions are typically at odds with the radio bounds. Note that these profiles are taken as representative cases of enhanced DM density distributions in the inner regions of the galaxy. However, the formation and evolution of a DM cusp at the galactic center is quite uncertain and other DM distributions falling between the two cases here considered, i.e. NFW and adiabatically contracted profiles, are possible (see Ref. [31,50]).

4.4 Signals at High Energy Colliders

Since the coupling of Z' to light SM fermions is suppressed by the small kinetic mixing factor, the best probe of the dark sector is through the top portal. In particular, the Z' can be produced by being radiated from top quarks, which have a large QCD production cross section at hadron colliders. In Fig. 13, we show the leading order cross section at LHC for $t\bar{t}Z'$ production, as calculated by CalcHep [56]. Depending on the masses and couplings, the Z' will predominantly decay into $t\bar{t}$, $\nu\bar{\nu}$, or into light fermions. Decays into top quarks lead to four-top events with a very large cross section compared to the SM four top rate, which can be visible through a same-sign dilepton signature [21] (see also [57] for studies

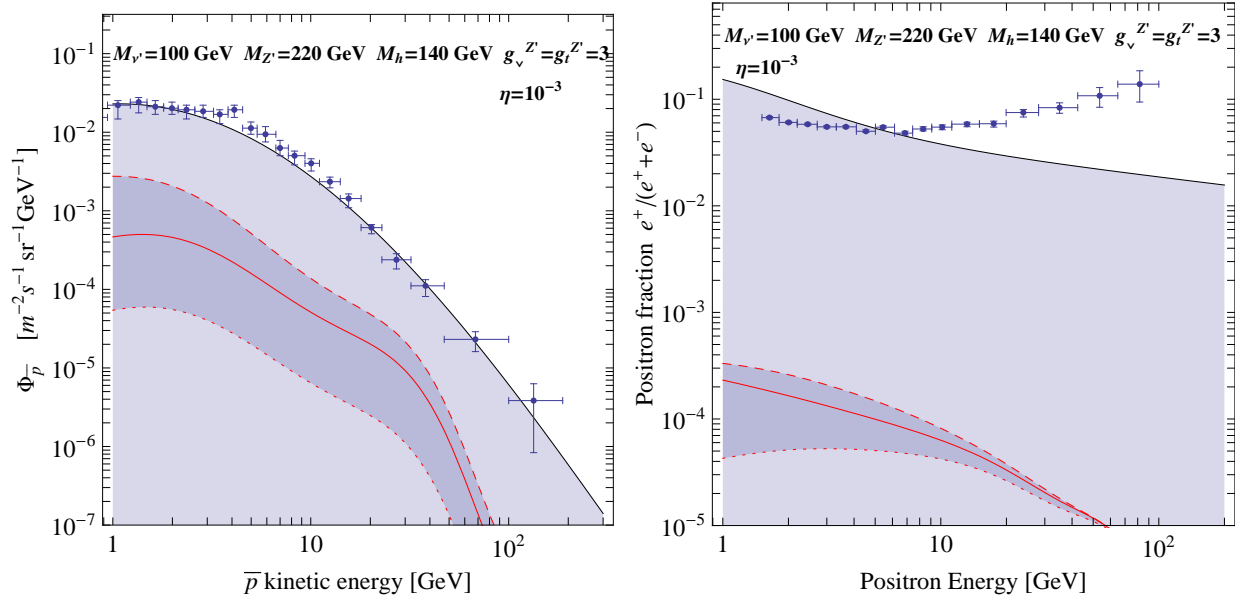


Figure 12: Left: antiproton flux measured by PAMELA [51]. Red lines show the contribution from ν annihilations for different sets of propagation parameters. In particular, dashed, solid and dotted red curves are respectively for the MAX, MED and MIN propagation models proposed in [52]. Right: positron fraction measured by PAMELA [53]. Dashed, solid and dotted red curves assume respectively the MAX, MED and MIN propagation models in [54]. In both panels, the black lines correspond to the expected astrophysical backgrounds. Antiproton fluxes, computed for a NFW profile, and background estimates are obtained as in [55].

of a $t\bar{t}WW$ final state). The right-handed nature of the Z' coupling to tops implies top polarization also provides an interesting observable. When the Z' decays into WIMPs, a $t\bar{t}$ +missing energy final state results, which presents a more challenging search at the LHC, but is definitely worth investigating. Work in these directions is in progress [58]. When the Z' is light, it may have large decays through its loop-induced interactions into γh or $b\bar{b}$. The γh decay provides a novel monoenergetic photon signature from $t\bar{t}Z' \rightarrow t\bar{t}h\gamma$. The decay into $b\bar{b}$ offers the possibility to reconstruct $M_{Z'}$ from the $b\bar{b}$ invariant mass of $t\bar{t}Z' \rightarrow t\bar{t}b\bar{b}$ events. Finally, the SM Higgs phenomenology at LHC could be changed due to the mixing with the Higgs responsible for the breaking of $U(1)'$. This was reviewed in [59]. Besides, in our model, the heavy Higgs could be produced by gluon fusion through a top loop and lead to interesting signatures in $t\bar{t}$ production.

5 Discussion and Outlook

In this article, we have shown that if dark matter has a large coupling to the top quark and suppressed couplings to light standard model degrees of freedom, as is expected when the DM dynamics are intimately linked with electroweak symmetry breaking, the Higgs could be produced copiously in the galaxy in association with a photon from dark matter

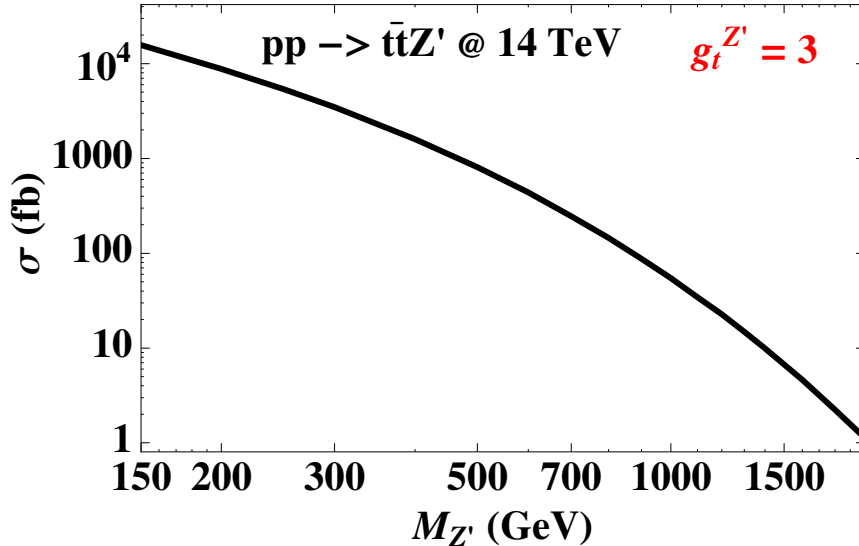


Figure 13: Production cross section of $t\bar{t}$ in association with a Z' . Above $M_{Z'} > 2m_t$, the branching fraction to top pairs can be large, resulting in a large $t\bar{t}t\bar{t}$ rate.

annihilations. The resulting photon spectrum possesses a line whose energy reflects the mass of the Higgs and of the WIMP and does not arise if DM is a scalar or a Majorana fermion, thus providing information about the WIMP spin and statistics. We have illustrated this phenomenon in a case in which the DM is composed of Dirac fermions annihilating via an s -channel massive vector resonance (Z') which couples strongly to the top quark, serving as a portal between dark matter and the Standard Model. Our setup arises naturally in models of “partial compositeness” in which the top quark acquires its large mass (after EWSB) through large mixing with composite states in a new strong sector, as in 4d duals to RS Models.

For couplings of order $\mathcal{O}(1)$, the correct dark matter abundance is reproduced from the standard thermal relic density calculation if the dark matter mass is of order the top mass, typically in the 100 GeV – 170 GeV mass range. In the limit of strong coupling, this feature is not strongly dependent on the mass of the Z' . This is a perfect mass range for searches with the Fermi LAT for gamma rays from WIMP annihilations, which we find has very good prospects for a discovery in the near future.

In part, gamma ray lines are particularly important because (unlike a typical model of WIMP DM, for which the photon continuum is usually much larger than the loop suppressed gamma ray lines), for $M_\nu < m_t$ (as is favored by the relic density in the strong coupling regime), the annihilation processes at the origin of the continuum photon emission are themselves a one loop process into $b\bar{b}$, enhancing the relative prominence of annihilation into γh and γZ . Even for $M_\nu \gtrsim m_t$, the continuum originating from annihilation into top quark pairs is rather soft, and the lines remain visible.

To illustrate Higgs production in space, we have worked with a simple representative

effective theory that captures properties of a general class of models with (some) simplifying assumptions. It would be interesting to refine our predictions for some particularly well-motivated models of electroweak symmetry breaking which fall into this category, such as technicolor models [60] or composite Higgs models [19]. In our effective theory construction, the WIMP effectively couples to the Higgs through the top quark, which may turn out to be ideal in terms of producing a γh line signal. Another interesting possibility would be to consider gauge-higgs unification models where the dark fermions couple directly to W gauge bosons [23,24]. Annihilation into γh would then be mediated by a box diagram involving W 's with couplings dictated by the gauge symmetry of the model. One could also easily imagine DM coupling directly to the Higgs through higher dimensional operators. In this work, we illustrated the γh signal from annihilating dark matter. However, also well-motivated are models where dark matter decay is induced by higher dimensional operators suppressed by a high scale close to the GUT scale, which, as mentioned in the introduction, can lead as well to gamma-ray line signatures. In fact, after this paper was posted, Ref. [61] appeared, which adequately completes our study by calculating the γh signal arising from decaying vector dark matter, leading naturally to an intense gamma-ray line.

As we explore the weak scale, we expect the dynamics of the electroweak breaking to be revealed. It may be that its secrets already shine down from the sky, produced by dark matter annihilation or decay.

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A A Simple 4d UV Completion

One particularly simple UV completion is to start with the Standard Model, treating all of its fields (including t_R) as uncharged under $U(1)'$. We include SM singlets ν_L and ν_R which are charged under $U(1)'$ to play the role of the WIMP, and in addition, a pair of fermions ψ_L and ψ_R , whose SM gauge quantum numbers are identical to t_R , but with equal charges under $U(1)'$. In this framework, the additional ingredients are vector-like, and thus the SM and mixed SM- $U(1)'$ anomalies are trivially absent. Depending on the charges of ν_L and ν_R ,

there may still be $U(1)^3$ anomalies, but these can be simply cancelled by adding SM gauge singlet fermions to the dark sector.

To realize coupling of the Z' to the top quark, we consider the gauge invariant masses and Yukawa couplings of the top- ψ sector,

$$yH\bar{Q}_3t_R + \mu\bar{\psi}_L\psi_R + Y\Phi\bar{\psi}_Lt_R \quad (37)$$

where Q_3 is the 3rd family quark doublet, H is the SM Higgs doublet, Φ is the Higgs field responsible for breaking $U(1)'$, y and Y are dimensionless couplings, and μ is a gauge-invariant mass term for ψ . In the regime $y\langle H \rangle \ll \mu, Y\langle \Phi \rangle$ and $\mu \sim Y\langle \Phi \rangle$, this system forms an ‘‘inverted top see-saw’’ in which the light eigenstate has mass $\sim y\langle H \rangle$ and is composed mostly of t_L and some fifty-fifty mixture of ψ_R and t_R , which we identify as the top quark. The heavier state is composed mostly of ψ_L and the orthogonal balanced mixture of t_R and ψ_R , whose mass $\sim \mu$ supplies the cut-off Λ in the effective theory.

B Coefficients and Vertex Factors for γZ and $\gamma Z'$

The coefficients (C_i) for the $Z'\gamma Z$ effective vertex are given by:

$$C_1 = \frac{-2(a_t g_V^t + v_t g_A^t)}{M_Z^2 - 4M_\nu^2} \left[M_Z^2 (B_0(4M_\nu^2; m_t^2, m_t^2) - B_0(M_Z^2; m_t^2, m_t^2)) \right] \\ + 2 \left[2m_t^2 (3a_t g_V^t - v_t g_A^t) C_0 + a_t g_V^t + v_t g_A^t \right] \quad (38)$$

$$C_2 = \frac{-2(a_t g_V^t + v_t g_A^t)}{M_Z^2 - 4M_\nu^2} \left[(4M_\nu^2 + M_Z^2) (B_0(4M_\nu^2; m_t^2, m_t^2) - B_0(M_Z^2; m_t^2, m_t^2)) \right. \\ \left. + 2(4M_\nu^2 - M_Z^2) (2m_t^2 C_0 + 1) \right] \quad (39)$$

$$C_3 = \frac{-4(a_t g_V^t + v_t g_A^t)}{(M_Z^2 - 4M_\nu^2)^2} \left[-M_Z^2 (B_0(4M_\nu^2; m_t^2, m_t^2) - B_0(M_Z^2; m_t^2, m_t^2)) \right. \\ \left. - (4M_\nu^2 - M_Z^2) (2m_t^2 C_0 + 1) \right] \quad (40)$$

$$C_4 = \frac{4(a_t g_V^t + v_t g_A^t)}{M_Z^2 - 4M_\nu^2} (B_0(4M_\nu^2; m_t^2, m_t^2) - B_0(M_Z^2; m_t^2, m_t^2)), \quad (41)$$

$$C_5 = \frac{4(a_t g_V^t + v_t g_A^t)}{(M_Z^2 - 4M_\nu^2)^2} \left[-M_Z^2 (B_0(4M_\nu^2; m_t^2, m_t^2) - B_0(M_Z^2; m_t^2, m_t^2)) \right. \\ \left. - (4M_\nu^2 - M_Z^2) (2m_t^2 C_0 + 1) \right], \quad (42)$$

where v_t and a_t are the SM vector and axial-vector couplings of a top quark to a Z boson, i.e.:

$$v_t = -1 + \frac{8}{3}s_W^2 \quad (43)$$

$$a_t = 1. \quad (44)$$

and, for the case where the Z' only couples to t_R , g_V^t and g_A^t are given by:

$$g_V^t = -g_A^t = \frac{g_t^{Z'}}{2}. \quad (45)$$

The corresponding coefficients for the $Z'\gamma Z'$ effective vertex can easily be obtained by replacing $a_t \rightarrow g_A^t$ and $v_t \rightarrow g_V^t$ in the above expressions.

Finally, the contribution to the γZ matrix-element-squared from the vertex factor $\mathcal{V}_{\gamma Z}^2$ is given in terms of the above C_i coefficients by:

$$\begin{aligned} \mathcal{V}_{\gamma Z}^2 = & |C_1|^2 \left(\frac{M_Z^2}{4} - 2M_\nu^2 + \frac{4M_\nu^4}{M_Z^2} \right) + M_Z^2 |C_2|^2 + |C_5|^2 \left(\frac{M_Z^6}{16} - M_\nu^2 M_Z^4 + 6M_\nu^4 M_Z^2 - 16M_\nu^6 + 16 \frac{M_\nu^8}{M_Z^2} \right) \\ & + (C_1 C_2^* + C_1^* C_2) \left(-\frac{M_Z^2}{2} + 2M_\nu^2 \right) + (C_1 C_5^* + C_1^* C_5) \left(-\frac{M_Z^4}{8} + \frac{3}{2} M_\nu^2 M_Z^2 - 6M_\nu^4 + \frac{8M_\nu^6}{M_Z^2} \right) \\ & + (C_2 C_5^* + C_2^* C_5) \left(\frac{M_Z^4}{4} - 2M_\nu^2 M_Z^2 + 4M_\nu^4 \right), \end{aligned} \quad (46)$$

where the corresponding factor for $\gamma Z'$ can be obtained from the above with the replacement $M_Z \rightarrow M_{Z'}$.

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