Decaying Hidden Dark Matter in Warped Compactification

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Abstract

The recent PAMELA and ATIC/Fermi/HESS experiments have observed an excess of electrons and positrons, but not anti-protons, in the high energy cosmic rays. To explain this result, we construct a decaying hidden dark matter model in string theory compactification that incorporates the following two ingredients, the hidden dark matter scenario in warped compactification and the phenomenological proposal of hidden light particles that decay to the Standard Model. In this model, on higher dimensional warped branes, various warped Kaluza-Klein particles and the zero-mode of gauge field play roles of the hidden dark matter or mediators to the Standard Model.
1 Introduction and model setup

More than eighty percent of the matter in the universe is dark matter, but its particle identities remains a mystery. Recent experiments may have observed the first non-gravitational signals of the dark matter. The PAMELA satellite [1] observed an excess of cosmic ray positron fraction at energies from 10 GeV up to at least 100 GeV, but importantly not the anti-proton excess in the same energy range. The ATIC balloon [2] found an excess of cosmic ray electrons above the astronomical background at energies 300-800 GeV with a peak at around 600 GeV. The recent Fermi satellite [3] also observed an excess of electrons in the similar energy range, but with a much softer spectrum comparing to ATIC. In addition, the HESS telescope [4] observed a steepening of electron spectrum above TeV. These electron or positron excess may be due to the annihilation or decay of the dark matter.

Given the high energy scale of the events, the conventional dark matter candidates such as the weakly interacting massive particles (WIMP) would have annihilated or decayed into both leptons and hadrons, including electrons/positrons and protons/anti-protons [5,6]. A possible phenomenology behind this anomaly is that the final decay to the SM model particles is due to a light particle [7–9]. The mass of this light particle is below about 1.8 GeV, so if decays on-shell the proton/anti-proton production is kinematically forbidden. In order for this to work, the annihilation channel or decay chain before that should be somewhat hidden from the Standard Model (SM) sector. This suggests the hidden dark matter scenarios [10–15]. In string theory compactification, sources carrying effective D3-brane charges create warped throats in six extra-dimensions [16–18]. During reheating, matter that is left in or tunnels to these throats can be trapped by gravitational potentials. If the SM is located somewhere else, these matter can become the hidden dark matter [10]. The main purpose of this paper is to construct models that incorporate the above two ingredients. We will combine the bottom-up and top-down approaches in constructing such models.

There are large number of ways a hidden sector may be implemented in terms of effective field theory. In string theory warped compactification, a natural candidate is the hidden throat. We will use such configurations to motivate hidden dark matter models that may or may not have simple 4d effective field theory descriptions. Various constraints, especially on the mediators between the hidden and SM sector, can arise from constructing models that have a reasonable UV completion. These may not be easy to envision from the pure bottom-up approach. On the other hand, experimental results play a crucial role in specifying our model configuration. This gives the motivation to explicitly construct some of the required components in string configuration.

In this paper we would like to construct a decaying hidden dark matter model in the
Figure 1: Configuration of the hidden throat. Crosses are the locations of the D3 or anti-D3-branes. Dashed lines indicate the higher dimensional branes; they wrap on the radial direction and the angular part $M_1$, and are embedded in $M_2$.

warped compactification, which meets the following model-building requirements:

1. The hidden dark matter is unstable and decays to the hidden light particles. This decay channel has a very long life time, such as of order $10^{26}$s.

2. The hidden light particles decay on-shell to the SM particles in an astrophysically short lifetime, before they can escape the galactic halos, which has a size 10 kpc. This requires the lifetime to be smaller than $10^{12}$s.\[1\]

3. The direct decay of the hidden dark matter into the SM particles such as proton/anti-protons and electron/positrons is much slower than the above two channels combined.

The following is our model configuration (Fig. 1). Consider a throat whose angular part is topologically a product, $M_1 \times M_2$. The radius for $M_1$ shrinks to zero, while $M_2$ remains finite at the bottom of the throat. We add higher dimensional spacetime-filling branes which wrap the radial direction and $M_1$. They can also wrap part of $M_1$, as long as they share an angular isometry in $M_1$. The additional, if any, dimensions of these branes are embedded in $M_2$. These higher dimensional branes extend outside the throat, and intersect with or become part of the SM branes in the bulk. Inside the throat, the gauge fields on the higher dimensional branes have warped Kaluza-Klein (WKK) excitations along the radial and $M_1$ directions; the $l$-th ($l > 0$) partial wave WKK particles become our hidden dark matter. Graviton also has WKK excitations and can become part of the hidden dark matter [10].\[2\]

As we will see, comparing to the gravitons, the zero-modes or the s-wave WKK modes of the

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1 I would like to thank Xuelei Chen and Xiaojun Bi for helpful discussions on this point.
2 If the SM is within the same throat, the stable gravity WKK particles [19] can also become viable dark matter candidates [10]. They belong to the conventional WIMP.
gauge fields are much stronger mediators to the SM. The minimum warped KK scale for this hidden throat is taken to be around TeV, which is an input from the ATIC/Fermi/HESS. At the bottom of the throat, there are spacetime-filling D3-branes or anti-D3-branes. Their locations can be separated in $M_2$. The open strings on these D3-branes provide light fields with masses hierarchically below TeV. This gives the required hidden light particle. We take its mass to be between 1.8 GeV and 1 MeV, so that it can decay to electron/positron, but not proton/anti-proton. This is an input from the PAMELA.

As mentioned, our model is motivated from both the existing flux compactification in type IIB string theory [16] and the experimental results. So on the one hand, the model will lead to a novel field theory of the hidden dark matter; on the other hand, not all components has been constructed explicitly in the current string theory literature. For example, before putting in D3-branes and wrapping higher dimensional branes, the required throat configuration is motivated from the Klebanov-Strassler (KS) throat [20], where the topology $M_1 \times M_2$ that we described is $S^2 \times S^3$. The higher dimensional branes can be D5 or D7-branes. However, the specific configuration of the D3-branes, the wrapped and warped D7-branes and their connection to the SM branes are proposed to satisfy the ATIC and PAMELA experiments. The number of the D7 and D3 branes should be small enough not to significantly affect the original throat.

When estimating lifetimes of various fields, we will concentrate on the most important factors regarding the above model building requirements. For the purpose of this paper, it suffices to ignore all the numerical pre-factors in our approximations. Many other models on the decaying dark matter [21] or the annihilating dark matter [22] have been studied in the context of the PAMELA and ATIC. Possible astrophysical explanations are studied in [23].

2 Gauge fields on warped brane

We consider the example of wrapping D7-branes on the whole $S^2$ of the KS throat. The $l$-th ($l > 0$) partial waves of the WKK modes have non-trivial angular dependence on $S^2$. When the size of $S^2$ shrinks to zero at the tip of the throat, their wavefunctions have to vanish to avoid the singularity, in contrast to the s-wave. So the wavefunctions do not overlap with the D3-branes located at the tip of the throat and these WKK modes are stable against decaying onto the D3-branes, in contrast to the s-wave. In other words, the presence of the D3-branes at the tip of the throat do not break the $S^2$ isometry associated with these partial waves, so we have conserved angular momenta. Once we are aware of this fact, we can ignore the very tip region of the throat and approximate the throat as the following simple form to study the characteristic behavior of the WKK wavefunction,

$$ds^2 = A^{-1/2}(-dt^2 + d\mathbf{x}^2) + A^{1/2}(dr^2 + r^2d\Omega_{S^2}^2 + r^2d\Omega_{S^3}^2),$$ (2.1)
where

\[ A = 1 + \frac{R^4}{r^4}. \]  

(2.2)

The \( r \) is the radial coordinate of the AdS space and \( R \) is the length scale of the warped space. The throat has a cut-off at \( r_0 \ll R \) and the minimum warp factor is \( h_0 \equiv r_0/R \).

The equation of motion for the Abelian gauge field \( B^{(8)}_\mu \) with a zero 8d mass on D7-branes is

\[ \partial_M \left( \sqrt{-G} G^{MP} G^{NQ} F^{(8)}_{PQ} \right) = 0 , \]  

(2.3)

where \( F^{(8)}_{PQ} = \partial_P B^{(8)}_Q - \partial_Q B^{(8)}_P \) is the field strength. We decompose \( B^{(8)}_\mu \equiv \tilde{B}^{(4)}_\mu (x) \varphi(r) \Phi(\Omega^{S^2}) \) and consider \( B^{(8)}_m = \text{const.} \). We use the Greek index \( \mu \) to label the 3+1 spacetime dimensions, and the Roman index \( m \) to label the extra-dimensions. The tilde in \( \tilde{B}^{(4)}_\mu \) indicates that it has not been canonically normalized. \( \Phi \) is the wavefunction on \( S^2 \), and the wavefunction on another angular direction is taken to be one. Choosing the Lorentz gauge \( \partial_\mu \tilde{B}^{(4)}_\mu = 0 \) and considering the plane-wave \( \tilde{B}^{(4)}_\mu \propto \epsilon_\mu e^{ipx} \), we get the differential equation for \( \varphi \) for the \( l \)-th partial wave in \( S^2 \),

\[ \frac{A^{1/4}}{r^2} \frac{d}{dr} \left( A^{-1/4} r^2 \frac{d}{dr} \varphi \right) - \frac{l(l+1)}{r^2} \varphi + m^2 A \varphi = 0 , \]  

(2.4)

where \( m^2 = -p^2 \) is the four-dimensional mass of the particle.

We are interested in the particles that are trapped inside the throat, so we impose the outgoing boundary condition. Solving (2.4), we get the following spectrum. It consists of a massless zero-mode, whose wavefunction on the D7-branes are constant, and a tower of WKK particles with masses quantized in unit of \( \sim h_0 R^{-1} \). At each level of the WKK, we have different partial waves labeled by \( l \), and their wavefunction in the radial direction behave as

\[ \varphi \propto r^{-b_l-1} \], \quad b_l = \sqrt{l^2 + l + 1} , \]  

(2.5)

for \( mR^2 \ll r \ll R; \varphi \propto r^{-l-1} \) for \( R \ll r \ll m^{-1} \). Comparing to the gravity modes, for the zero-mode, the world-volume of D7-branes on which the gauge fields propagate is only a much smaller subspace of the whole bulk; for the WKK modes, the damping of the wavefunction is much slower than that of the gravity which is for example \( r^{-4} \) for the s-wave [10]. These facts will become important later. The zero-mode disappears if the gauge field acquires a 8d mass.
3 Dark matter decay within hidden throat

In a realistic compactification, the angular isometries of the throat will always be broken, for example at least by the presence of the SM branes outside the throat. Let us consider the effect of a stack of $\tilde{N}$ extra D3-branes located at $r = y_0$. Instead of (2.2), the warped geometry is now a superposition of that of the throat and branes,

$$A = 1 + \frac{R^4}{r^4} + \frac{\tilde{N} R^4}{N |r - y_0|^4},$$

$$= 1 + \frac{R^4}{r^4} + \frac{\tilde{N} R^4}{N y_0^4} \left( 1 + \frac{4r}{y_0} \cos \theta + \cdots \right), \quad (3.6)$$

where we have expanded around $r = 0$ for $r \ll y_0$, $\theta$ is the angle between $r$ and $y_0$, $N$ is the charge of the hidden throat.

The angular dependence in (3.6) introduces a mixing between the $l$-th ($l \neq 0$) and $s$-wave WKK mode in the following otherwise vanishing kinetic term,

$$\sim \frac{m_8^4}{g_s} \int d^8x \sqrt{-G^{(8)}} F^{(8)}_{MNPQ} G^{(8)}_{MNPQ} G^{(8)}_{NPQRS}, \quad (3.7)$$

where one of $F^{(8)}_{MNP}$ is for the $l$-th partial wave and another the $s$-wave. Integrating the extra dimensions using the wavefunction (2.5) and the deformed metric (3.6), and normalizing the four-dimensional fields $\tilde{B}^{(4)}_\mu$ using (3.7) for the $s$-wave and $l$-th partial wave respectively, we get the mixing

$$\epsilon \sim \tilde{h}_0^5 \frac{\tilde{N}}{N} \left( \frac{R}{y_0} \right)^5. \quad (3.8)$$

Due to the wavefunction (2.5), the integration is dominated by the contribution from the tip of the throat for $l \geq 2$. The parameter (3.8) is predominant over the other model-dependent parameters that we will assume on the hidden particle model on D3-branes.

On the D3-branes at the tip of the throat, there are light particles. The 8d $s$-wave WKK gauge field intersects with the D3-branes and induces a 4d gauge field through a mixing strength $\alpha_1$, which is determined by the string coupling $g_s$ between the open strings on D7-branes and D3-branes. The 4d gauge field couples to other light particles, for example, the fermion-anti-fermion $\psi / \bar{\psi}$ with a coupling constant $e_1$. Neither $\alpha_1$ nor $e_1$ depends on the warp factor. So the coupling of the $s$-wave $B^{(8)}_\mu$ to $\psi / \bar{\psi}$ is given by

$$e_1 \alpha_1 \int d^4x \bar{\psi} \gamma^\mu B^{(8)}_\mu \psi,$$

where $B^{(8)}_\mu$ is evaluated at the location of the D3-branes.
Working with the canonically normalized fields $B^{(4)}_\mu$, using (3.8) and (3.9), we get the decay rate of the $l$-th WKK to two $\psi$'s,

$$\Gamma_{l\to s\to 2\psi} \sim h_0^{10} \left( \frac{R}{y_0} \right)^{10} \left( \frac{N}{N_3} \right)^{10} \epsilon_1^2 \alpha_1^2 m_{WKK}. \quad (3.10)$$

We have used the relation $m_{WKK} \sim h_0 R^{-1}$ and $R \sim (g_s N)^{1/4} m_s^{-1}$. The fermions can quickly cascade into many lighter particles on D3-branes, such as a neutral boson $\chi$ that is stable before decaying outside the throat.

When we integrate the extra-dimensions in (3.7), the integrand is proportional to $r^{-b_l+1}$. For $l \geq 2$, the main contribution comes from the tip region. This is what we used above. For $l = 1$, the integration gets most contribution toward the UV region, and we integrate it up to the location of the extra branes, $y_0$. The mixing becomes larger. In fact, for low $l$, there is another channel that can give a competitive decay rate. That is, the $l$-th WKK can also decay through mixing with the zero-mode gauge field. Because the zero-mode has a constant wavefunction, in the integration over the extra dimensions, the integrand is less peaked at the tip of the throat. This integration gets most contribution near the UV region for $l \leq 3$. The zero-mode gauge field couples to $\psi/\bar{\psi}$ through the same term (3.9). The decay rate through this channel is

$$\Gamma_{l\to 0\to 2\psi} \sim h_0^{2b_l+2} \left( \frac{R}{y_0} \right)^{2b_l+2} \left( \frac{R}{L_b} \right)^8 \left( \frac{N}{N_3} \right)^8 \epsilon_1^2 \alpha_1^2 m_{WKK}, \quad l \leq 3. \quad (3.11)$$

Although the power on the warp factor is smaller than (3.10), (3.11) gets an extra suppression from the D7-brane extra-dimensional volume $L_b^4$ because $L_b$ is larger than the size of the throat $R$. As we will discuss in Sec. 5, this is a channel we want to avoid. For $l \geq 4$, a calculation similar to (3.10) shows that the decay rate through this channel is smaller than (3.10) by a factor of $(R/L_b)^8$, so it is always highly suppressed.

To give a numerical example, suppose $e_1 \sim 0.1$, $\alpha_1 \sim 0.1$, $N \sim 10$, $N \sim 100$ and $y_0 \sim R$. The mass of the WKK is determined by the ATIC/Fermi/HESS to be around TeV. In order for the lifetime of the dark matter to be $10^{26}$s, from (3.10) we get the warp factor of the hidden throat, $h_0 \sim 10^{-4.5}$. It is exponentially small, as naturally realized in the flux compactification.

It is clear that in this scenario the decay rate is highly dependent of the environment, such as the warp factor of the hidden throat and the location of the extra branes. This is also seen for the gravity WKK modes as the conventional WIMP [24–26]. For the gravity WKK modes, it is possible to systematically classify different perturbations that break the

\[3\text{The lifetime of the dark matter may be shorter if the light particles which eventually decay to electrons/positrons have a small production branching ratio here.} \]
angular isometries [24–27]. It may be done similarly here, although the coefficients of the classifications still need to be determined by a detailed configuration. We have used a simple configuration to convey the following main message regarding the requirement 1 in Sec. 1. The isometry breaking, that is present in the bulk or UV side of the throat, is separated from the IR tip by the gravitational potential of the warped geometry, so the decay rate is typically suppressed by a large power of the warp factor. Hence the lifetime of such a TeV scale hidden dark matter is very long. However the precise decay rate and the energy scale of the cosmic rays are typically not the specific predictions of general hidden dark matter models [10]. In fact, it is natural in this scenario that there are multiple peaks of cosmic rays at various energy scales and associated with different dark matter lifetimes. We are just starting to observe one of them. We now turn to the next two requirements concerning the various communications between the hidden throat and the SM.

4 Decay of hidden light particle to SM

The decay of the hidden light particle $\chi$ to the SM electron/positron can be mediated by the gauge fields on the D7-branes because they extend outside the throat and intersect with the SM located at a distance $D$. The zero-mode and $s$-wave are most important for this purpose, since higher partial waves damp faster. A Feynman diagram of such a decay is given in Fig. 2.

We first consider the zero-mode. The particle in the loop is a hidden fermion that couples to the hidden D3-brane gauge field. As mentioned in Sec. 3, this induces a coupling to the D7-brane gauge field, and in turn induces a coupling to the SM electro-magnetic gauge field. The first vertex is given by a Yukawa coupling $\lambda_1$. Both the second and third vertex are given by the term (3.9), at different locations and with different couplings $e_{1,2}$, $\alpha_{1,2}$ and particle contents $\psi, \psi_e$. In terms of the normalized 4d field $B^{(4)}_{\mu}$, both couplings receive a suppression factor from the D7-brane volume, $(R/L_b)^2 N^{-1/2}$. This is the most important parameter for this process. Comparing to the $s$-wave mediation that we will study shortly, there is no warp factor here because the wavefunction of the zero-mode is constant on D7-branes. Comparing

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Figure 2: Feynman diagram for the decay of the light particle $\chi$ to electron/positron.
to the graviton zero-mode, which effectively gives a coupling $m_{\text{WKK}}/M_{\text{Pl}} \sim h_0(R/L)^3N^{-1}$ ($L$ is the size of the bulk), the coupling here is much stronger for the following reasons. First, the D7-branes are embedded in the bulk, so its size is typically smaller, $L_b < L$; second, the D7-brane has fewer dimensions to integrate over; third, the couplings $e_{1,2}$, $\alpha_{1,2}$ are already dimensionless so it is not affected by warping.

The decay rate of $\chi$ through this channel is

$$\Gamma_{\chi \rightarrow 0 \rightarrow 2e} \sim \frac{h_0^4}{L_b^2} \frac{\Lambda_{1}^2 e_{1}^2 e_{2}^2 \alpha_{1}^2 \alpha_{2}^2}{m_{\chi}^4} \frac{E_0}{\gamma}.$$  \hspace{1cm} (4.12)

The $\Lambda$ is the momentum cutoff in the loop of $\psi$, as an example we approximate it to be $m_{\text{WKK}} \sim \text{TeV}$ which is the warped KK scale at the bottom of the throat. We assume $m_{\chi} \gg 2m_e$. The energy scale in the center-of-mass frame is $E_0 \sim m_{\chi}$, and $\gamma \sim m_{\text{WKK}}/m_{\chi}$ is the Lorentz factor for $\chi$. The denominator $m_{\chi}^4$ comes because the zero-mode mediator is massless. So if there are several stable light particles as $\chi$, the lightest one is most favorable to decay in this channel.

For the $s$-wave mediation, the most important factor is the suppression factor in the third vertex, $h_0^2(R/D)$ for $D > R$ and $h_0^2(R/D)^2$ for $D < R$, due to the WKK damping outside the throat. The decay rate through this channel is

$$\Gamma_{\chi \rightarrow s \rightarrow 2e} \sim h_0^4 \frac{R}{D} \frac{1}{N^2} \lambda_{1}^2 \epsilon_{1}^2 \epsilon_{2}^2 \alpha_{1}^2 \alpha_{2}^2 \frac{\Lambda_{1}^4}{m_{\text{WKK}}^4} \frac{E_0}{\gamma}.$$  \hspace{1cm} (4.13)

The factor $(R/D)^2$ becomes $(R/D)^4$ for $D < R$. The heavier $\chi$ is more favorable to decay in this channel.

Let us use the same numerical example in Sec. 3 and assume $\lambda_1 \sim e_2 \sim \alpha_2 \sim 0.1$. The $R/L_b$ can be as small as $R/L \sim (Nh_0^{-1}m_{\text{WKK}}/M_{\text{Pl}})^{1/3} \sim 10^{-3}$, so we use $R/L_b \sim 10^{-2}$ for example. The decay rate $\Gamma_{\chi \rightarrow 0 \rightarrow 2e}$ then ranges from $10^3$ to $10^9s^{-1}$, for $m_{\chi}$ from GeV to MeV. The decay mediated by the $s$-wave is much slower, $\Gamma_{\chi \rightarrow s \rightarrow 2e} \sim 10^{-11}$ to $10^{-17}s^{-1}$. For some parameter space this lifetime is still short enough to satisfy our requirement. This channel becomes important in case the zero-mode gets lifted. If we instead use the $s$-wave gravity WKK \cite{10} to mediate the decay, we will have an extra factor of $h_0^4$. This makes the lifetime of $\chi$ easily exceed the age of the universe\cite{5}.

5 Direct decay of dark matter to SM

A direct decay of the hidden dark matter to the SM should produce both electrons/positrons and protons/anti-protons, among many other particles. A rate larger than (3.10) would

\footnote{The case where both the hidden dark matter and hidden light particle have lifetimes longer than the age of the universe is a viable case, but the decay products are no longer observable since they escape the galactic halos.}
generically contradict the PAMELA experiment. There are a couple of ways a direct decay can happen.

First, since the wavefunction of the WKK dark matter itself has a damping tail outside the throat, the direct decay can be mediated by the term \( B_{\mu}^{(8)} \), but now we use the \( l \)-th partial wave in \( B_{\mu}^{(8)} \) and evaluate it at the location of the SM branes. The decay rate is

\[
\Gamma_{l \to 2e} \sim h_0^{2l+2} \left( \frac{R}{D} \right)^{2l+2} \frac{1}{N} e_2^2 \alpha_2^2 m_{\text{WKK}}. \tag{5.14}
\]

The factor \( (R/D)^{2l+2} \) becomes \( (R/D)^{2l+2} \) for \( D < R \). The decay rate to the proton/anti-proton or other hadrons is the same except that the couplings \( e_2 \) and \( \alpha_2 \) may be different.

Eq. (5.14) decreases drastically as \( l \) increases, because larger angular momentum introduces higher effective potential. Comparing with (3.10), for \( l \geq 4 \), (5.14) has at least additional powers of warp factor and is highly suppressed. So all requirements in Sec. 1 are naturally satisfied. For \( l \leq 3 \), (5.14) has a smaller power on \( h_0 \). In order to suppress (5.14), one can decrease \( y_0 \) to increase the isometry-breaking and increase \( D \) to move away the SM branes. Neither case affects the \( \chi \) decay rate (4.12). (But both affect (4.13).) In this case, we are introducing isometry breaking sources besides the SM branes. For the most difficult case, \( l = 1 \), using (4.12) it is enough to have, for example, \( y_0 \sim 10^{-3} R \), \( D \sim 10^2 R \) and \( h_0 \sim 10^{-7.5} \).

Second, the direct decay can be mediated by mixing with virtual particles. The channel through a virtual \( s \)-wave WKK is doubly suppressed by both mixing and tunneling. For the virtual zero-mode, we need to suppress the decay rate (3.11) relative to (3.10). This is because the zero-mode gauge field couples similarly to the D3-branes both inside and outside the throat through the term (3.9), so this channel mediates the hadron production. Again, this suppression is naturally achieved for \( l \geq 4 \). For \( l \leq 3 \), especially \( l = 1 \), we need to significantly increase the D7-brane volume suppression. In our previous example, \( R/L_b \) can be as small as \( 10^{-3} \), which is enough for the suppression to work for all \( l \leq 3 \).

Alternatively, the absence of the first few low-\( l \) (\( l > 0 \)) partial waves can be achieved by some discrete symmetries in the angular directions. For example, a discrete symmetry on the \( S^2 \) azimuthal angle under \( \varphi \to \pi/2 + \varphi \) makes the partial wave start from \( l = 4 \).

6 Summary and discussion

In this paper we have constructed a decaying hidden dark matter model in warped compactification. Generally speaking, in this class of models, there is an angular isometry shared by the hidden throat and wrapped higher dimensional branes, and this isometry is not broken by the presence of the D3-branes (or anti-D3-branes) at the tip of the throat. This requirement is necessary only if the hidden dark matter is the WKK modes discussed in Sec. 2 & 3. It can be dropped if the hidden dark matter is something else in the hidden particle sector (e.g. analogous
Figure 3: The decay channels of the WKK dark matter. The solid lines are the dominant decay channels. The dashed line is a much slower channel. The dotted lines are highly suppressed. Curved lines indicate the virtual particles involved in the Feynman diagrams.

a mass hierarchy between the WKK modes (of effective fields from open or closed strings) associated with this isometry and the light particles on the D3-branes, which serve as the hidden dark matter and the hidden light particles, respectively. The communication between the hidden throat and the SM sector is mediated by the zero-mode or the s-wave WKK mode of the gauge field (or other similar fields) on the higher dimensional branes.

We studied an example of such a model and found the following main features (Fig. 3).

1. The objects that break the isometry associated with the WKK modes are separated by the gravitational potential of the throat from the WKK particles. So the hidden dark matter, either gauge field or gravity WKK particles, have very long lifetime.

2. The decay of light particles to SM electrons/positrons can be mediated by the zero-mode or the lowest s-wave WKK mode of the gauge field on the higher dimensional branes. The zero-mode is efficient (especially for interaction terms with dimensionless coupling constants), because its wavefunction does not damp as WKK, and the volume suppression is much weaker than the gravity zero-mode. This mediation mechanism is very important for the communication between the hidden sector and SM. In case to the lightest supersymmetric particles for SM). In the latter case, the only requirement for the higher dimensional branes is that they intersect (or nearly intersect) both the SM and hidden particle sectors. The mediation between the two sectors is provided by the zero-mode gauge field on the higher dimensional branes.

By construction, the mediator here is electro-magnetically neutral and conserves the SM lepton and hadron number. So the stable SM particles do not decay through this mediator.
the zero-mode gets lifted, the s-wave mediation may also be sufficient (although much weaker), because its wavefunction damps much slower than the gravity WKK.

3. The direct decay of the dark matter to the SM particles is highly suppressed. This is because the direct tunneling is suppressed by the effective potentials caused by the warped space and WKK’s angular momenta; while the channel through the virtual gauge field zero-mode (or s-wave WKK) is suppressed by both the small mixing and the volume of the higher dimensional branes (or the tunneling).

Many aspects are worth to be further studied. In a realistic warped compactification, there are other components and fields in the hidden throat besides what we have considered. It will be interesting to explore their roles in this model. It is also interesting to construct the required brane configuration more explicitly in string theory. At least, any SM particles lighter than the electron, i.e. neutrino and photon, are kinematically allowed in the decay, but can have different branching ratios; on the hidden D3-branes, there can also naturally be massless particles. The associated observational effects such as the high energy gamma ray or neutrino emissions is an important issue. It is also important to study possible collider physics signals.

Acknowledgments

I would like to thank Xiaojun Bi, Xuelei Chen, Alejandro Jenkins, Igor Klebanov and Henry Tye for helpful discussions and comments. This work is supported by the US Department of Energy under cooperative research agreement DEFG02-05ER41360.
References

[1] O. Adriani et al., arXiv:0810.4995 [astro-ph].

[2] J. Chang et al., Nature 456, 362 (2008).

[3] A. A. Abdo et al. [The Fermi LAT Collaboration], arXiv:0905.0025 [astro-ph.HE].

[4] H. E. S. Aharonian, arXiv:0905.0105 [astro-ph.HE].

[5] M. Cirelli, M. Kadastik, M. Raidal and A. Strumia, arXiv:0809.2409 [hep-ph].

[6] F. Donato, D. Maurin, P. Brun, T. Delahaye and P. Salati, arXiv:0810.5292 [astro-ph].

[7] D. P. Finkbeiner and N. Weiner, Phys. Rev. D 76, 083519 (2007) arXiv:astro-ph/0702587.

[8] I. Cholis, L. Goodenough and N. Weiner, arXiv:0802.2922 [astro-ph].

[9] N. Arkani-Hamed, D. P. Finkbeiner, T. Slatyer and N. Weiner, arXiv:0810.0713 [hep-ph].

[10] X. Chen and S. H. Tye, JCAP 0606, 011 (2006) arXiv:hep-th/0602136.

[11] T. Kikuchi and N. Okada, Phys. Lett. B 665, 186 (2008) arXiv:0711.1506 [hep-ph].

[12] B. v. Harling and A. Hebecker, JHEP 0805, 031 (2008) arXiv:0801.4015 [hep-ph].

[13] Y. Gong and X. Chen, Eur. Phys. J. C 57, 785 (2008) arXiv:0803.3223 [astro-ph].

[14] J. L. Feng and J. Kumar, Phys. Rev. Lett. 101, 231301 (2008) arXiv:0803.4196 [hep-ph].

[15] J. L. Feng, H. Tu and H. B. Yu, JCAP 0810, 043 (2008) arXiv:0808.2318 [hep-ph].

[16] S. B. Giddings, S. Kachru and J. Polchinski, Phys. Rev. D 66, 106006 (2002) arXiv:hep-th/0105097.

[17] L. Randall and R. Sundrum, Phys. Rev. Lett. 83, 3370 (1999) arXiv:hep-ph/9905221.

[18] H. L. Verlinde, Nucl. Phys. B 580, 264 (2000) arXiv:hep-th/9906182.

[19] L. Kofman and P. Yi, Phys. Rev. D 72, 106001 (2005) arXiv:hep-th/0507257.

[20] I. R. Klebanov and M. J. Strassler, JHEP 0008, 052 (2000) arXiv:hep-th/0007191.
[21] C. R. Chen, F. Takahashi and T. T. Yanagida, Phys. Lett. B 671, 71 (2009) [arXiv:0809.0792 [hep-ph]].

P. f. Yin, Q. Yuan, J. Liu, J. Zhang, X. j. Bi and S. h. Zhu, arXiv:0811.0176 [hep-ph].

K. Ishiwata, S. Matsumoto and T. Moroi, arXiv:0811.0250 [hep-ph].

C. R. Chen, M. M. Nojiri, F. Takahashi and T. T. Yanagida, arXiv:0811.3357 [astro-ph].

E. Nardi, F. Sannino and A. Strumia, arXiv:0811.4153 [hep-ph].

M. Pospelov and M. Trott, arXiv:0812.0432 [hep-ph].

A. Arvanitaki, S. Dimopoulos, S. Dubovsky, P. W. Graham, R. Harnik and S. Rajendran, arXiv:0812.2075 [hep-ph].

K. Hamaguchi, S. Shirai and T. T. Yanagida, arXiv:0812.2374 [hep-ph].

F. Takahashi and E. Komatsu, arXiv:0901.1915 [astro-ph].

K. Hamaguchi, F. Takahashi and T. T. Yanagida, arXiv:0901.2168 [hep-ph].

C. H. Chen, C. Q. Geng and D. V. Zhuridov, arXiv:0901.2681 [hep-ph].

B. Kyae, arXiv:0902.0071 [hep-ph].

[22] J. H. Huh, J. E. Kim and B. Kyae, arXiv:0809.2601 [hep-ph].

N. Arkani-Hamed and N. Weiner, JHEP 0812, 104 (2008) [arXiv:0810.0714 [hep-ph]].

A. E. Nelson and C. Spitzer, arXiv:0810.5167 [hep-ph].

I. Cholis, D. P. Finkbeiner, L. Goodenough and N. Weiner, arXiv:0810.5344 [astro-ph].

Y. Nomura and J. Thaler, arXiv:0810.5397 [hep-ph].

R. Harnik and G. D. Kribs, arXiv:0810.5557 [hep-ph].

D. Feldman, Z. Liu and P. Nath, arXiv:0810.5762 [hep-ph].

Y. Bai and Z. Han, arXiv:0811.0387 [hep-ph].

P. J. Fox and E. Poppitz, arXiv:0811.0399 [hep-ph].

I. Cholis, G. Dobler, D. P. Finkbeiner, L. Goodenough and N. Weiner, arXiv:0811.3641 [astro-ph].

G. Bertone, M. Cirelli, A. Strumia and M. Taoso, arXiv:0811.3744 [astro-ph].

K. M. Zurek, arXiv:0811.4429 [hep-ph].

M. Taoso, S. Ando, G. Bertone and S. Profumo, arXiv:0811.4493 [astro-ph].

M. Ibe, H. Murayama and T. T. Yanagida, arXiv:0812.0072 [hep-ph].
J. Hisano, M. Kawasaki, K. Kohri and K. Nakayama, arXiv:0812.0219 [hep-ph].
E. J. Chun and J. C. Park, arXiv:0812.0308 [hep-ph].
M. Lattanzi and J. I. Silk, arXiv:0812.0360 [astro-ph].
J. Zhang, X. J. Bi, J. Liu, S. M. Liu, P. f. Yin, Q. Yuan and S. H. Zhu, arXiv:0812.0522 [astro-ph].
J. D. March-Russell and S. M. West, arXiv:0812.0559 [astro-ph].
R. Allahverdi, B. Dutta, K. Richardson-McDaniel and Y. Santoso, arXiv:0812.2196 [hep-ph].
D. Hooper, A. Stebbins and K. M. Zurek, arXiv:0812.3202 [hep-ph].
K. J. Bae, J. H. Huh, J. E. Kim, B. Kyae and R. D. Viollier, arXiv:0812.3511 [hep-ph].
L. Bergstrom, G. Bertone, T. Bringmann, J. Edsjo and M. Taoso, arXiv:0812.3895 [astro-ph].
C. R. Chen, K. Hamaguchi, M. M. Nojiri, F. Takahashi and S. Torii, arXiv:0812.4200 [astro-ph].
P. Grajek, G. Kane, D. Phalen, A. Pierce and S. Watson, arXiv:0812.4555 [hep-ph].
X. J. Bi, P. H. Gu, T. Li and X. Zhang, arXiv:0901.0176 [hep-ph].
M. Baumgart, C. Cheung, J. T. Ruderman, L. T. Wang and I. Yavin, arXiv:0901.0283 [hep-ph].
S. C. Park and J. Shu, arXiv:0901.0720 [hep-ph].
I. Gogoladze, R. Khalid, Q. Shafi and H. Yuksel, arXiv:0901.0923 [hep-ph].
S. Khalil, H. S. Lee and E. Ma, arXiv:0901.0981 [hep-ph].
Q. H. Cao, E. Ma and G. Shaughnessy, arXiv:0901.1334 [hep-ph].
W. L. Guo and Y. L. Wu, arXiv:0901.1450 [hep-ph].
P. Meade, M. Papucci and T. Volansky, arXiv:0901.2925 [hep-ph].
J. Mardon, Y. Nomura, D. Stolarski and J. Thaler, arXiv:0901.2926 [hep-ph].
D. J. Phalen, A. Pierce and N. Weiner, arXiv:0901.3165 [hep-ph].
J. Hisano, M. Kawasaki, K. Kohri, T. Moroi and K. Nakayama, arXiv:0901.3582 [hep-ph].
F. Chen, J. M. Cline and A. R. Frey, arXiv:0901.4327 [hep-ph].
[23] D. Hooper, P. Blasi and P. D. Serpico, JCAP 0901, 025 (2009) [arXiv:0810.1527 [astro-ph]].

H. Yuksel, M. D. Kistler and T. Stanev, arXiv:0810.2784 [astro-ph].

S. Profumo, arXiv:0812.4457 [astro-ph].

K. Ioka, arXiv:0812.4851 [astro-ph].

R. Brandenberger, Y. F. Cai, W. Xue and X. Zhang, arXiv:0901.3474 [hep-ph].

[24] A. Berndsen, J. M. Cline and H. Stoica, Phys. Rev. D 77, 123522 (2008) [arXiv:0710.1299 [hep-th]].

[25] J. F. Dufaux, L. Kofman and M. Peloso, Phys. Rev. D 78, 023520 (2008) [arXiv:0802.2958 [hep-th]].

[26] J. M. Cline, R. J. Danos and A. R. Frey, work in progress.

[27] A. Ceresole, G. Dall'Agata and R. D'Auria, JHEP 9911, 009 (1999) [arXiv:hep-th/9907216].