



Gamma rays and positrons from a decaying hidden gauge boson

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ABSTRACT

We study a scenario that a hidden gauge boson constitutes the dominant component of dark matter and decays into the standard model particles through a gauge kinetic mixing. Interestingly, gamma rays and positrons produced from the decay of hidden gauge boson can explain both the EGRET excess of diffuse gamma rays and the HEAT anomaly in the positron fraction. The spectra of the gamma rays and the positrons have distinctive features; the absence of line emission of the gamma ray and a sharp peak in the positron fraction. Such features may be observed by the FGST and PAMELA satellites.

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1. Introduction

The concept of symmetry has been the guiding principle in modern physics. The structure of the standard model (SM) is dictated by $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetries. The electroweak symmetry, $SU(2)_L \times U(1)_Y$, is spontaneously broken by a non-vanishing vacuum expectation value (vev) of the Higgs boson, and massive W and Z bosons are generated. It is quite natural in the string landscape that there are many other gauge symmetries as well as discrete ones realized in nature, and some of the gauge symmetries may be spontaneously broken, leading to massive gauge bosons, as in the SM.

Suppose that some of the hidden gauge bosons are 'odd' under a Z_2 parity, whereas all the SM particles are 'even'. If the parity is exact, and if the hidden gauge symmetries are spontaneously broken, the lightest parity-odd gauge boson is stable and can be a good candidate for cold dark matter. Some recent examples are the T-odd $U(1)$ gauge boson in the little Higgs model with T-parity [1–4] and the KK photon in the minimal universal extra dimension model [5]. However, it is not known yet which of unbroken or (explicitly or spontaneously) broken discrete symmetries are more common in the string landscape. If a discrete symmetry breaking is a general phenomenon, we may expect that the dark matter is not absolutely stable and ultimately decays into the SM particles.

In this Letter we consider a simplest case that there exist a hidden $U(1)'$ gauge symmetry which is spontaneously broken and a parity under which only the hidden gauge boson changes its sign. We assume that the parity is broken by a kinetic mixing between the $U(1)'$ and the SM $U(1)_Y$ gauge symmetries [6,7]. As a result,

the hidden gauge boson decays into the SM particles through the parity violating interactions induced by the kinetic mixing. If such a violation is so tiny that the lifetime of the hidden gauge boson is much longer than the age of our universe, the hidden gauge boson can be the dominant component of dark matter. Furthermore, the subsequent decays of the SM particles will form continuous spectra of the gamma rays and the positrons in the high-energy cosmic ray. With an appropriate amount of the parity violation, we see that those gamma rays and positrons could be the source of the excesses of the gamma ray observed by Energetic Gamma Ray Experiment Telescope (EGRET) [8,9] and the positron flux observed by High Energy Antimatter Telescope (HEAT) [10], MASS [11] and AMS [12] experiments.

Recently, the gravitino dark matter scenario was extensively studied in the framework of supersymmetry with R -parity violation [13,14]. The decay of the gravitino into the gamma rays was discussed in Refs. [13,14]. Furthermore, if the gravitino is heavier than W boson, it decays predominantly into a W or Z boson and a lepton. With the gravitino mass of $\mathcal{O}(100)$ GeV and the lifetime of $\mathcal{O}(10^{26})$ s, the gravitino decay can simultaneously explain the EGRET anomaly in the extragalactic diffuse gamma ray background and the HEAT excess in the positron fraction [15–18]. Our decaying hidden gauge boson has therefore many parallels with the gravitino with R -parity violation. However, in our model, the decay branching ratio of the hidden gauge boson to W boson is highly suppressed, which is significantly different from the gravitino case.

This Letter is organized as follows. In Section 2, we present the effective Lagrangian and Feynman rules used in our calculations. In Section 3, we calculate the spectra of gamma ray and positron flux from the decay of the hidden gauge boson and compare them with the observed data. Section 4 is devoted to discussions and conclusions.

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2. Framework

We consider a hidden Abelian gauge symmetry $U(1)'$ and the associated gauge boson B'_μ , and introduce a hidden parity under which B'_μ transforms as

$$B'_\mu \rightarrow -B'_\mu. \quad (1)$$

The $U(1)'$ symmetry is assumed to be spontaneously broken, so that the gauge boson B'_μ has a non-vanishing mass, m .

We assume that the low energy effective theory can be written in terms of the SM particles and the hidden gauge boson B'_μ , and that all the SM particles are neutral under both the $U(1)'$ and the hidden parity. We would like to introduce a tiny parity violating interaction. Among many possibilities, we will focus on a kinetic mixing term between $U(1)'$ and $U(1)_Y$, since the kinetic mixing has the lowest dimension.¹ The low energy effective Lagrangian can be written as

$$\mathcal{L} \supset -\frac{1}{4}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4}B'_{\mu\nu}B'^{\mu\nu} + \frac{\epsilon}{2}B_{\mu\nu}B'^{\mu\nu} + \frac{1}{2}m^2B'_\mu B'^\mu, \quad (2)$$

where ϵ ($\ll 1$) is the coefficient of the hidden parity violation term, and B_μ is the gauge boson of $U(1)_Y$, $B_{\mu\nu}$ and $B'_{\mu\nu}$ are the field strength of $U(1)_Y$ and $U(1)'$, respectively: $B_{\mu\nu}^{(i)} = \partial_\mu B_\nu^{(i)} - \partial_\nu B_\mu^{(i)}$. In the following discussion we neglect $\mathcal{O}(\epsilon^2)$ terms, since ϵ must be extremely small as $\mathcal{O}(10^{-26})$ for the hidden gauge boson to become dark matter. We will also give a possible origin of such tiny kinetic mixing later.

We can remove the kinetic mixing and bring the kinetic terms into canonical form by redefining the gauge fields as [19]

$$\begin{aligned} \tilde{B} &= B - \epsilon B', \\ \tilde{B}' &= B'. \end{aligned} \quad (3)$$

Note that we omit the Lorentz index μ hereafter. Taking account of the electroweak symmetry breaking, the mass terms of neutral gauge bosons are

$$\begin{aligned} \mathcal{L}_{\text{mass}}^{\text{Gauge}} &= \frac{v^2}{8} \begin{pmatrix} W^3 & \tilde{B} & \tilde{B}' \end{pmatrix} \\ &\times \begin{pmatrix} g^2 & -gg' & -gg'\epsilon \\ -gg' & g'^2 & g'^2\epsilon \\ -gg'\epsilon & g'^2\epsilon & 4m^2/v^2 \end{pmatrix} \begin{pmatrix} W^3 \\ \tilde{B} \\ \tilde{B}' \end{pmatrix}, \end{aligned} \quad (4)$$

where W^3 is the neutral gauge field of $SU(2)_L$, v is the vev of the doublet Higgs field H , g and g' are weak and hypercharge coupling constants, respectively. After diagonalizing the mass matrix, we can express the gauge eigenstates W^3 , B and B' in terms of the mass eigenstates Z , A , and A' as

$$\begin{aligned} W^3 &= c_W Z + s_W A - c_W s_W \frac{m_Z^2}{m_{A'}^2 - m_Z^2} \epsilon A', \\ B &= -s_W Z + c_W A + \frac{m_{A'}^2 - c_W^2 m_Z^2}{m_{A'}^2 - m_Z^2} \epsilon A', \\ B' &= s_W \frac{m_Z^2}{m_{A'}^2 - m_Z^2} \epsilon Z + A' \end{aligned} \quad (5)$$

with

¹ It is also possible to write down non-renormalizable parity violating interactions suppressed by a large mass scale. For the magnitudes of kinetic mixing in our work, however, we can safely neglect such interactions if they are suppressed by the Planck scale.

$$\begin{aligned} c_W &\equiv \cos \theta_W = g / \sqrt{g^2 + g'^2}, & s_W &\equiv \sin \theta_W = \sqrt{1 - c_W^2}, \\ m_Z^2 &= \frac{1}{4}(g^2 + g'^2)v^2, & m_{A'}^2 &= m^2, \end{aligned} \quad (6)$$

where $m_{A'}$ and m_Z are the masses of A' and Z , respectively. One can easily see that A and Z are reduced to the ordinary photon and Z boson in the limit of $\epsilon \rightarrow 0$. Since interested values of ϵ are extremely small, the kinetic mixing hardly affects any SM predictions of the electroweak measurements.

The A' interacts with the SM particles through the mixings shown in Eq. (5). The Feynman rules can be derived in a straightforward way by expanding the SM Lagrangian with respect to ϵ [20,21], and the results are

$$\begin{aligned} \bar{u}A'_\mu u: & -i \frac{g'\epsilon}{2(m_{A'}^2 - m_Z^2)} \gamma_\mu \left(-\frac{4}{3}c_W^2 m_Z^2 + \frac{5}{6}m_{A'}^2 + \frac{1}{2}m_{A'}^2 \gamma_5 \right), \\ \bar{d}A'_\mu d: & -i \frac{g'\epsilon}{2(m_{A'}^2 - m_Z^2)} \gamma_\mu \left(\frac{2}{3}c_W^2 m_Z^2 - \frac{1}{6}m_{A'}^2 - \frac{1}{2}m_{A'}^2 \gamma_5 \right), \\ \bar{\nu}A'_\mu \nu: & i \frac{g'\epsilon m_{A'}^2}{4(m_{A'}^2 - m_Z^2)} \gamma_\mu (1 - \gamma_5), \\ \bar{e}A'_\mu e: & -i \frac{g'\epsilon}{4(m_{A'}^2 - m_Z^2)} \gamma_\mu (4c_W^2 m_Z^2 - 3m_{A'}^2 - m_{A'}^2 \gamma_5), \\ WW A': & -\frac{s_W m_Z^2 \epsilon}{m_{A'}^2 - m_Z^2} g_{WWZ}^{\text{SM}}, \end{aligned} \quad (7)$$

where u , d , e and ν represent all the up-type quarks, down-type quarks, charged leptons and neutrinos, respectively, and g_{WWZ}^{SM} is the coupling of WWZ in the SM.

The hidden gauge boson A' decays into the SM particles through the above interactions. Fig. 1 shows the decay branching ratios and the lifetime of A' as a function of $m_{A'}$. The down-type quark decay modes of A' , $A' \rightarrow d\bar{d}$ ($d = d, s, b$), dominate over the other modes for $m_{A'} \simeq 100$ GeV, whereas they decrease quickly as $m_{A'}$ increases. For $m_{A'} \gtrsim 120$ GeV, the up-type quark decay modes, $A' \rightarrow u\bar{u}$ ($u = u, c$), become the largest ones followed by charged lepton decay modes, $A' \rightarrow \ell^+\ell^-$ ($\ell = e, \mu, \tau$). This behavior can be understood easily as follows. The A' becomes more like the Z boson as $m_{A'}$ approaches m_Z , and therefore, the partial decay width of $A' \rightarrow f\bar{f}$ is proportional to $g_V^2 + g_A^2$, where $g_{V(A)}$ is the vector (axial) coupling strength of the neutral weak interaction. On the other hand, as $m_{A'}$ becomes much heavier than m_Z , A' tends to not feel the electroweak symmetry breaking, so the branching ratios are insensitive to $m_{A'}$. For $m_{A'} \gtrsim 350$ GeV, the decay mode of $A' \rightarrow t\bar{t}$ is allowed. The following discussions on the gamma rays and the positrons are not significantly modified even for $m_{A'} \gtrsim 350$ GeV.

3. Gamma ray and positron spectra

The hidden gauge boson A' decays into a pair of SM fermions and W bosons if kinematically allowed, as we have seen in the previous section. Due to the subsequent QCD hadronization processes, a bunch of hadrons are produced, and in particular, continuum spectra for the photon and the positron are formed. Throughout this Letter we assume that the A' constitutes the dominant component of dark matter. Then the produced photons and positrons may be observed in the high-energy cosmic rays. In this section we estimate the fluxes of the gamma rays and the positrons produced from the A' decay, and see how they may account for the observed excesses.

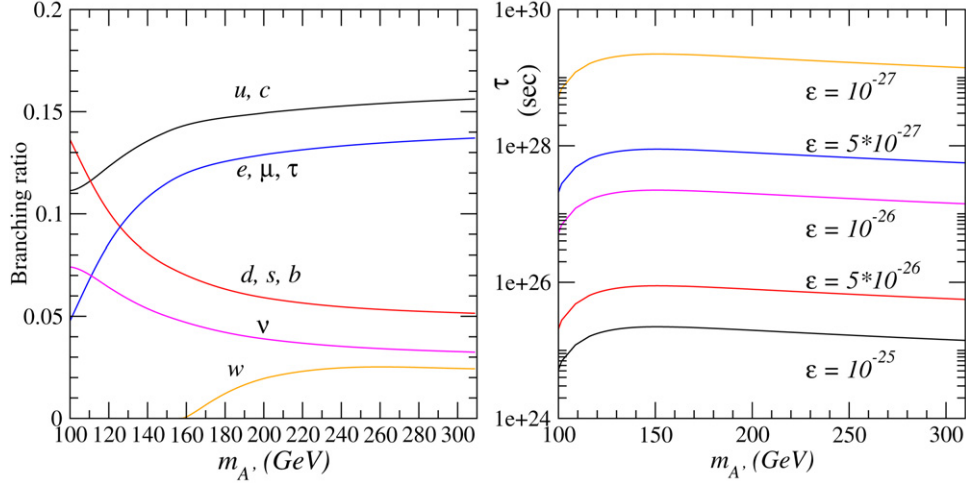


Fig. 1. Decay branching ratios and lifetime of A' , as a function of $m_{A'}$.

3.1. Gamma-ray flux

The gamma-ray energy spectrum is characterized by dN_γ/dE , the number of photons having energy between E and $E + dE$, produced from the decay of one A' gauge boson. The main contribution to the continuous spectrum of γ arises from the π^0 generated in the QCD hadronization process. To estimate the spectrum, we use the PYTHIA [22] Monte Carlo program with the branching ratios shown in Fig. 1. We also include the real gamma-ray emission, known as internal bremsstrahlung, from the charged particles from the A' direct decay,² whose contributions to dN_γ/dE become important compared to that from π^0 when $E \rightarrow m_{A'}/2$. However, the features of the gamma-ray flux which we discuss below will not change even though the internal bremsstrahlung effects are not included. The energy spectra dN_γ/dE for $m_{A'} = 100$ and 300 GeV are shown in Fig. 2. It is worth noting that there is no line emission of the gamma rays from the decay of A' , which is present in the case of the gravitino dark matter.

There are galactic and extragalactic contributions from the decay of A' to the observed gamma ray flux. The flux of the gamma ray from the extragalactic origin is estimated as [16,18]

$$\left[E^2 \frac{dJ_\gamma}{dE} \right]_{\text{eg}} = \frac{E^2 c \Omega_{A'} \rho_c}{4\pi m_{A'} \tau_{A'} H_0 \Omega_M^{1/2}} \int_1^{y_{\text{eq}}} dy \frac{dN_\gamma}{d(yE)} \frac{y^{-3/2}}{\sqrt{1 + \frac{\Omega_\Lambda}{\Omega_M} y^{-3}}}, \quad (8)$$

where c is the speed of light; $\Omega_{A'}$, Ω_M and Ω_Λ are the density parameters of A' , matter (including both baryons and dark matter) and the cosmological constant, respectively; ρ_c is the critical density; $\tau_{A'}$ is the lifetime of A' ; H_0 is the Hubble parameter at the present time; $y \equiv 1 + z$, where z is the redshift, and y_{eq} denotes a value of y at the matter-radiation equality. For the numerical results, we use [24]

$$\Omega_{A'} h^2 = 0.1099, \quad \Omega_M h^2 = 0.1326, \quad \Omega_\Lambda = 0.742, \\ \rho_c = 1.0537 \times 10^{-5} \text{ GeV/cm}^3. \quad (9)$$

On the other hand, the gamma ray flux from the decay of A' in the Milky Way halo is

$$\left[E^2 \frac{dJ_\gamma}{dE} \right]_{\text{halo}} = \frac{E^2}{4\pi m_{A'} \tau_{A'}} \frac{dN_\gamma}{dE} \left\langle \int_{\text{los}} \rho_{\text{halo}}(\vec{\ell}) d\vec{\ell} \right\rangle, \quad (10)$$

² We are grateful to John Beacom [23] for bringing up this fact to us.

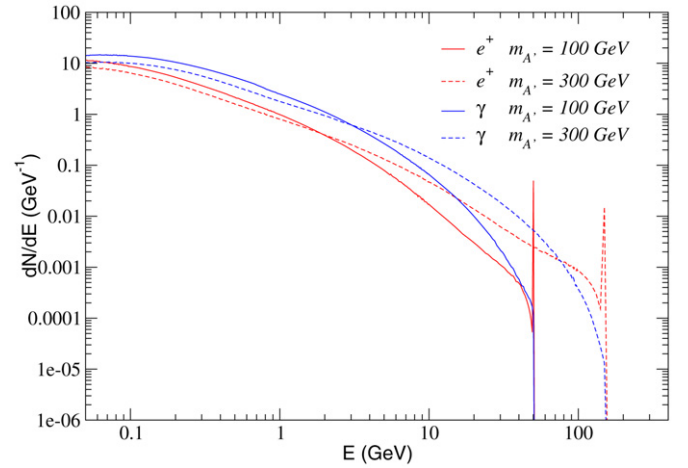


Fig. 2. Energy spectra of γ and e^+ generated from the decay of A' .

where ρ_{halo} is the density profile of dark matter in the Milky Way, $\langle \int_{\text{los}} \rho_{\text{halo}}(\vec{\ell}) d\vec{\ell} \rangle$ is the average of the integration along the line of sight (los). We adopt the Navarro-Frenk-White (NFW) halo profile [25]

$$\rho(r) = \frac{\rho_0}{(r/r_c)(1 + r/r_c)^2} \quad (11)$$

in our calculation, where r is the distance from the center of Milky Way, $r_c = 20$ kpc, and ρ_0 is set in such a way that the dark matter density in the solar system satisfies $\rho(r_\odot) = 0.30 \text{ GeV/cm}^3$ [26] with $r_\odot = 8.5$ kpc being the distance from the Sun to the Galactic Center.

In order to compare the EGRET results with the above flux from A' decay, we integrate over the whole sky except for the zone of the Galactic plane (i.e. the region with the galactic latitudes $|b| < 10^\circ$). For the background, we use a power-law form adopted in Ref. [18]

$$\left[E^2 \frac{dJ_\gamma}{dE} \right]_{\text{bg}} \simeq 5.18 \times 10^{-7} E^{-0.499} \text{ GeV cm}^{-2} \text{ sr}^{-1} \text{ s}^{-1}, \quad (12)$$

where E is in units of GeV. Fig. 3 shows our numerical results. Recall that there are only two parameters in our model, i.e. the mass and lifetime of A' , and the peak position only depends on the mass while the deviation from background is sensitive to both. The signal of $m_{A'} = 100$ GeV (the red line) peaks at around $E = 5$ GeV region, which is consistent very well with the observed data. For

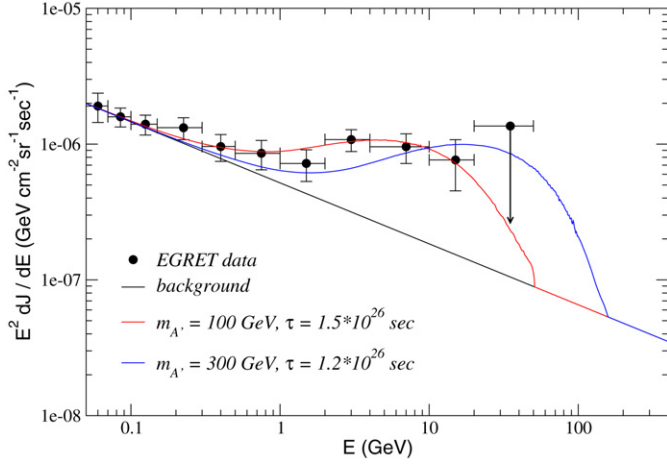


Fig. 3. γ ray flux predicted from decay of A' , shown together with the EGRET data [9]. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

$m_{A'} = 300$ GeV (the blue line), our prediction can still account for the observed excess, although the fit is not so good in some energy region. However, given the large errors in the observed data, it may be premature to extract any sensible constraint on the model parameters. With more precise data from the FGST experiment, we should be able to have more information about the A' mass and lifetime, if the excess is indeed from A' dark matter decay. As mentioned before, there is no line emission of the gamma rays in our model, because the production of a pair of on-shell γ from A' is forbidden [27,28]. Therefore, there is no secondary peak around the high end of signal region, which is a characteristic difference between our prediction and the gravitino case.

3.2. Positron fraction

After being produced from the A' decay, the positron will propagate in the magnetic field of the Milky Way. The typical gyroradius is much smaller than the size of the galaxy, and the positron will propagate along the magnetic field. However, since the magnetic fields are tangled, the motion of the positron can be described by a diffusion equation. Neglecting the convection and annihilation in the disk, the steady state solution must satisfy [17]

$$\nabla \cdot [K(E, \vec{r}) \nabla f_{e^+}] + \frac{\partial}{\partial E} [b(E, \vec{r}) f_{e^+}] + Q(E, \vec{r}) = 0, \quad (13)$$

where f_{e^+} is the number density of e^+ per unit kinetic energy, $K(E, \vec{r})$ is the diffusion coefficient, $b(E, \vec{r})$ is the rate of energy loss and $Q(E, \vec{r})$ is the source of producing e^+ from A' decay. In our case,

$$Q(E, \vec{r}) = \frac{\rho(\vec{r})}{m_{A'} \tau_{A'}} \frac{dN_{e^+}}{dE},$$

where dN_{e^+}/dE is the energy spectrum of e^+ from A' decay obtained by using PYTHIA [22] (see Fig. 2). We notice that because of the decay channel of $A' \rightarrow e^+e^-$, there is a sharp peak at $E = m_{A'}/2$.

The solution of Eq. (13) in the solar system can be expressed as

$$f_{e^+}(E) = \frac{1}{m_{A'} \tau_{A'}} \int_0^{E_{\text{Max}}} dE' G(E, E') \frac{dN_{e^+}}{dE'}, \quad (14)$$

where $E_{\text{Max}} = m_{A'}/2$, and $G(E, E')$ is approximately given by [17]

$$G(E, E') \simeq \frac{10^{16}}{E^2} e^{a+b(E^{\delta-1}-E'^{\delta-1})} \theta(E' - E) \text{ s/cm}^3, \quad (15)$$

where E is in units of GeV, δ is related to the properties of the interstellar medium and can be determined mainly from the ratio of boron to carbon (B/C) [29]. We adopt parameters, $\delta = 0.55$, $a = -0.9716$ and $b = -10.012$ [17], that are consistent with the B/C value and produce the minimum flux of positrons. Finally, the flux of e^+ is given by

$$\Phi_{e^+}^{\text{prim}}(E) = \frac{c}{4\pi} f_{e^+} = \frac{c}{4\pi m_{A'} \tau_{A'}} \int_0^{m_{A'}/2} dE' G(E, E') \frac{dN_{e^+}}{dE'}. \quad (16)$$

In addition to e^+ flux from dark matter decay, there exists a secondary e^+ flux from interactions between cosmic rays and nuclei in the interstellar medium. The positron flux is considered to be suffered from the solar modulation, especially for the energy below 10 GeV. If the solar modulation effect is independent of the charge-sign, one can cancel the effect by measuring the positron fraction,

$$\frac{\Phi_{e^+}}{\Phi_{e^+} + \Phi_{e^-}}. \quad (17)$$

Indeed, most of experiments measured the fraction of positron flux. To estimate the positron fraction, it is necessary to include the e^- flux. We use the approximations of the e^- and e^+ background fluxes [30,31]

$$\begin{aligned} \Phi_{e^-}^{\text{prim}}(E) &= \frac{0.16E^{-1.1}}{1 + 11E^{0.9} + 3.2E^{2.15}} \text{ GeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \\ \Phi_{e^-}^{\text{sec}}(E) &= \frac{0.7E^{0.7}}{1 + 110E^{1.5} + 600E^{2.9} + 580E^{4.2}} \text{ GeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \\ \Phi_{e^+}^{\text{sec}}(E) &= \frac{4.5E^{0.7}}{1 + 650E^{2.3} + 1500E^{4.2}} \text{ GeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \end{aligned} \quad (18)$$

where E is in units of GeV. Therefore, the fraction of e^+ flux is

$$\frac{\Phi_{e^+}^{\text{prim}} + \Phi_{e^+}^{\text{sec}}}{\Phi_{e^+}^{\text{prim}} + \Phi_{e^+}^{\text{sec}} + k\Phi_{e^-}^{\text{prim}} + \Phi_{e^-}^{\text{sec}}}, \quad (19)$$

where k is a free parameter which is used to fit the data when no primary source of e^+ flux exists [31,32]. Note also that the primary flux of e^- , $\Phi_{e^-}^{\text{prim}}$, in the denominator of Eq. (19) should include the contributions from dark matter A' decay as well. Our numerical results for $m_{A'} = 100$ GeV (the magenta line) and 300 GeV (the green line) are shown in Fig. 4. The prediction of our model is consistent with the observed excess quite well, and the positron fraction starts increasing around $E \simeq 20$ GeV and $E \simeq 10$ GeV for $m_{A'} = 100$ GeV and $m_{A'} = 300$ GeV, respectively. Another key feature of the signal prediction is that the curve drops off sharply at $E = m_{A'}/2$ mainly due to the contribution of e^+ from $A' \rightarrow e^+e^-$ decay channel, i.e. the peak seen in Fig. 2. These characteristics can be checked by the upcoming PAMELA data.

4. Discussion and conclusions

Let us here discuss briefly how such a small kinetic mixing in Eq. (2) can arise. If there is an unbroken Z_2 parity symmetry under which the hidden gauge boson flips its sign, the kinetic mixing would be forbidden. This parity may arise from a more fundamental symmetry or it may be an accidental one. In order to have a non-vanishing kinetic mixing, we need to break the parity by a small amount and transmit the breaking to the kinetic mixing.³ To this end we introduce messenger fields ξ_1 and ξ_2 that have both

³ With the presence of the parity symmetry, a small breaking is natural in the sense of 't Hooft.

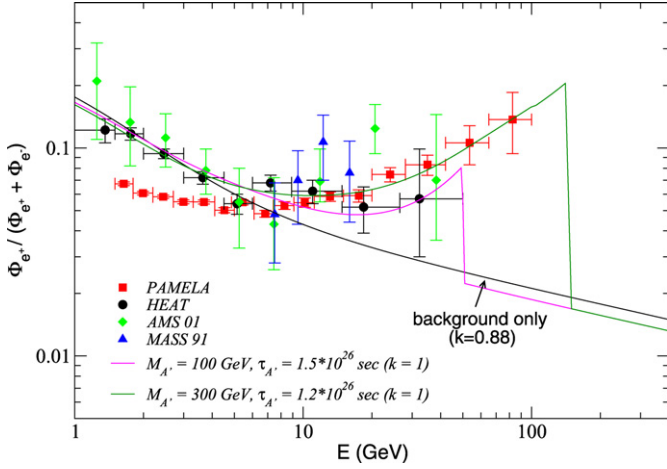


Fig. 4. Fraction of e^+ flux from A' decay, shown together with experimental data [10–12,33]. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article Letter.)

$U(1)_Y$ and $U(1)'$ charges as $(1, 1)$ and $(1, -1)$, respectively. Under the parity, they transform as

$$\xi_1 \leftrightarrow \xi_2. \quad (20)$$

Note that the charge assignment is consistent with the parity transformation (see the discussion below Eq. (23)). We assume that the messenger mass scale M_m is very large, e.g. the grand unified theory (GUT) scale. Suppose that the parity is broken in the messenger sector in such a way that ξ_1 and ξ_2 obtain different masses. Let us denote the mass difference by ΔM_m , which parametrizes the amount of the parity violation. Integrating out the messengers, we are then left with a small kinetic mixing, ϵ , given by

$$\epsilon \sim \frac{g_h g'}{16\pi^2} \frac{\Delta M_m^2}{M_m^2}, \quad (21)$$

where g_h is the coupling constant of the hidden $U(1)$ gauge symmetry. For $\Delta M_m \sim \mathcal{O}(\text{TeV})$ and $M_m \sim 10^{15}$ GeV, we obtain the tiny kinetic mixing of the right magnitude of $\mathcal{O}(10^{-26})$.

It is tempting to identify the origin of the parity violation with the spontaneous breaking of the $U(1)'$. As an illustration, we will present a toy model below. As we will see later, there are some dangerous couplings which may spoil the stability of the A' in this model. However, these problems could be solved by embedding the model into a theory with supersymmetry or an extra dimension(s).

Suppose that all the matter fields in the hidden sector are neutral under the parity. Then, the gauge field B'_μ cannot have any interactions with those hidden matter fields, since they are forbidden by the parity. Let us introduce two scalar fields, ϕ_1 and ϕ_2 , which transform under the parity as

$$\phi_1 \leftrightarrow \phi_2, \quad (22)$$

and we assume that ϕ_1 and ϕ_2 are neutral under the SM gauge symmetries. Then the following interactions are allowed:

$$\mathcal{L} \supset i g_h B'_\mu (\phi_1^* \partial^\mu \phi_1 - \text{h.c.}) - i g_h B'_\mu (\phi_2^* \partial^\mu \phi_2 - \text{h.c.}) \quad (23)$$

The parity is spontaneously broken if one of the two scalars develops a non-vanishing vev. To this end, we consider the following potential:

$$\lambda (|\phi_1|^2 - v_h^2)^2 + \lambda (|\phi_2|^2 - v_h^2)^2 + 2\kappa |\phi_1|^2 |\phi_2|^2, \quad (24)$$

where v_h , λ and κ are real and positive. For $\kappa > \lambda$, there are four distinct vacua, $(\phi_1, \phi_2) = (\pm v_h, 0)$ and $(0, \pm v_h)$. We take one possibility of them, $(\phi_1, \phi_2) = (v_h, 0)$, as an example. Therefore, the hidden $U(1)$ gauge symmetry is spontaneously broken by

the vev of ϕ_1 , and the associated gauge boson acquires a mass, $m = \sqrt{2} g_h v_h$, by eating the imaginary component of ϕ_1 .

In order to transmit the parity violation, we introduce couplings between $\phi_{1,2}$ and the messenger fields,

$$-\mathcal{L} \supset M_m^2 (|\xi_1|^2 + |\xi_2|^2) + \kappa' (|\phi_1|^2 |\xi_1|^2 + |\phi_2|^2 |\xi_2|^2), \quad (25)$$

where M_m is the messenger mass, and κ' is a real and positive constant. After the $U(1)'$ is spontaneously broken, masses of ξ_1 and ξ_2 are slightly different: $m_{\xi_1}^2 = M_m^2 + \kappa' v_h^2$ and $m_{\xi_2}^2 = M_m^2$. After integrating out these heavy messengers, we obtain the kinetic mixing in Eq. (2) with ϵ given by

$$\epsilon \sim \frac{g_h g'}{16\pi^2} \frac{\kappa' v_h^2}{M_m^2}.$$

For $g_h \sim \kappa' \sim \mathcal{O}(1)$, $v_h \sim 1$ TeV, and $M_m \sim 10^{15}$ GeV, we obtain $\epsilon \sim 10^{-26}$. We can therefore realize more or less the correct magnitude of ϵ needed to account for the excesses of the gamma rays and the positrons in this toy model.

In the above toy model, we have assumed that the SM particles couple to the hidden sector only through the messenger fields. It is also possible to introduce direct interactions between $\phi_{1,2}$ and the visible sector. For instance, we can couple them to the SM fermions as

$$\mathcal{L} \supset \frac{\alpha}{M_p^2} (|\phi_1|^2 + |\phi_2|^2) \bar{f}_L f_R H + \text{h.c.}, \quad (26)$$

where $\alpha \sim \mathcal{O}(1)$, $M_p \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck scale, and $f_{R(L)}$ is the right-(left)-handed fermion and H is the Higgs field. Although this interaction does not lead to the gauge kinetic mixing, it induces the decay of the hidden gauge boson into a fermion pair after breaking the parity. However, the decay branching ratio through such interaction is negligible, compared to that through the kinetic mixing with $\epsilon \sim 10^{-26}$.

More dangerous direct couplings are those between $\phi_{1,2}$ and the Higgs field, $(|\phi_1|^2 + |\phi_2|^2)|H|^2$. If there exist such interactions, the A' could decay into the Higgs bosons immediately. If we extend the model into a supersymmetric one, the direct couplings to the Higgs bosons can be suppressed with the aid of the R -symmetry. In a theory with an extra dimension, it is also possible to suppress the direct couplings with appropriate configurations of the branes (e.g. the visible particles on one brane, while the hidden particles on the other).

It is also known that many $U(1)$ symmetries appear in the string theory, and the kinetic mixing can similarly arise in the low energy theory by integrating out heavy string states that have charges of two $U(1)$ gauge symmetries. It has been extensively studied how large the kinetic mixing can be in e.g. Ref. [34] (see also Ref. [35] and references therein). For instance, in a warped background geometry, we can have an exponentially small kinetic mixing [35].

Let us also have some comments on the GUT. So far we have assumed the existence of the kinetic mixing between $U(1)_Y$ and the hidden $U(1)'$. If one of the $U(1)$ gauge symmetries actually sits within an unbroken non-Abelian gauge symmetry, such kinetic mixing is not allowed. This does not necessarily mean that kinetic mixings are incompatible with GUT, because the GUT gauge group is spontaneously broken. Indeed, kinetic mixings between the $U(1)'$ and the GUT gauge fields can arise below the GUT scale by picking up non-vanishing vevs of the Higgs bosons responsible for the GUT breaking.

Let us also discuss how the hidden gauge boson could be generated in the early universe to account for the observed abundance of dark matter. For simplicity, we neglect a numerical coefficient of order unity in the following discussion. In the presence of the

messenger fields, there appears at one-loop level a following interaction between the $U(1)'$ and $U(1)_Y$ gauge fields:

$$\mathcal{L} \supset \frac{g_h^2 g'^2}{16\pi^2 M_m^4} (B_{\alpha\beta} B^{\alpha\beta}) (B'_{\gamma\delta} B'^{\gamma\delta}). \quad (27)$$

The hidden gauge boson A' will be produced through the above interaction most efficiently at the reheating. The A' abundance is roughly estimated to be

$$\frac{n_{A'}}{s} \sim 10^{-12} g_h^4 \left(\frac{T_R}{3 \times 10^{13} \text{ GeV}} \right)^7 \left(\frac{M_m}{10^{15} \text{ GeV}} \right)^{-8}, \quad (28)$$

where s and T_R are the entropy density and the reheating temperature, respectively. For the mass $m_{A'} \sim \mathcal{O}(100)$ GeV, a right abundance of A' can be generated from the above interaction for $T_R \sim 10^{13}$ GeV. Also the A' can be non-thermally produced by the inflaton decay [36–38].

In the above model, the parity is spontaneously broken by the vev of ϕ_1 . If the breaking occurs after inflation, domain walls connecting two of the four vacua in Eq. (24) will be formed [39–41], which can be the cosmological disaster. There are several means to get around this problem, and one of which is to introduce a tiny explicit breaking of the parity symmetry. Therefore the domain walls are not stable, and eventually annihilate after collisions [41–43]. Another solution is to assume that the breaking occurs before inflation. The last one is to assume that the initial positions of the scalars ϕ_1 and ϕ_2 are deviated from the origin. In the last case, the domain walls, if formed, will be annihilated eventually.⁴

In this Letter we have considered a possibility that a hidden gauge boson A' , which constitutes the dominant component of dark matter, decays into the SM particles through the kinetic mixing term that breaks the Z_2 parity symmetry. As a result, the branching ratios are solely determined by the mass of the hidden gauge boson. Continuum spectra of photons and positrons are generated from $A' \rightarrow \ell^+ \ell^-$ ($\ell = e, \mu, \tau$), and from the decays of hadrons, mainly $\pi^{0,\pm}$, produced in the subsequent QCD hadronization process.⁵ If the mass of A' is about $\mathcal{O}(100)$ GeV and its lifetime is of order $\mathcal{O}(10^{26})$ seconds, those gamma rays and positrons from A' decay may account for the observed excesses in the extra galactic diffuse gamma ray flux and the positron fraction. Interestingly, in our model, the spectra of the gamma rays and the positrons have distinctive features: the absence of line emission of the gamma ray and a sharp peak in the positron fraction. Such features may be observed by the FGST and PAMELA satellites.

Note added

Very recently the PAMELA group reported a steep rise in the positron fraction [33], which is nicely explained by our model.

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References

- [1] N. Arkani-Hamed, A.G. Cohen, E. Katz, A.E. Nelson, JHEP 0207 (2002) 034, hep-ph/0206021.
- [2] H.-C. Cheng, I. Low, JHEP 0309 (2003) 051, hep-ph/0308199.
- [3] H.-C. Cheng, I. Low, JHEP 0408 (2004) 061, hep-ph/0405243.
- [4] C. Csaki, J. Heinonen, M. Perelstein, C. Spethmann, arXiv: 0804.0622.
- [5] T. Appelquist, H.-C. Cheng, B.A. Dobrescu, Phys. Rev. D 64 (2001) 035002, hep-ph/0012100.
- [6] B. Holdom, Phys. Lett. B 166 (1986) 196.
- [7] R. Foot, X.-G. He, Phys. Lett. B 267 (1991) 509.
- [8] P. Sreekumar, et al., EGRET Collaboration, Astrophys. J. 494 (1998) 523, astro-ph/9709257.
- [9] A.W. Strong, I.V. Moskalenko, O. Reimer, Astrophys. J. 613 (2004) 956, astro-ph/0405441.
- [10] S.W. Barwick, et al., HEAT Collaboration, Astrophys. J. 482 (1997) L191, astro-ph/9703192.
- [11] C. Grimani, et al., Astron. Astrophys. 392 (2002) 287.
- [12] M. Aguilar, et al., AMS-01 Collaboration, Phys. Lett. B 646 (2007) 145, astro-ph/0703154.
- [13] F. Takayama, M. Yamaguchi, Phys. Lett. B 485 (2000) 388, hep-ph/0005214.
- [14] W. Buchmuller, L. Covi, K. Hamaguchi, A. Ibarra, T. Yanagida, JHEP 0703 (2007) 037, hep-ph/0702184.
- [15] G. Bertone, W. Buchmuller, L. Covi, A. Ibarra, JCAP 0711 (2007) 003, arXiv: 0709.2299.
- [16] A. Ibarra, D. Tran, Phys. Rev. Lett. 100 (2008) 061301, arXiv: 0709.4593.
- [17] A. Ibarra, D. Tran, JCAP 0807 (2008) 002, arXiv: 0804.4596.
- [18] K. Ishiwata, S. Matsumoto, T. Moroi, arXiv: 0805.1133, 2008.
- [19] T. Fukuda, T. Yanagida, M. Yonezawa, Prog. Theor. Phys. 51 (1974) 1406.
- [20] W.-F. Chang, J.N. Ng, J.M.S. Wu, Phys. Rev. D 74 (2006) 095005, hep-ph/0608068.
- [21] S. Gopalakrishna, S. Jung, J.D. Wells, arXiv: 0801.3456.
- [22] T. Sjostrand, S. Mrenna, P. Skands, JHEP 0605 (2006) 026, hep-ph/0603175.
- [23] J.F. Beacom, N.F. Bell, G. Bertone, Phys. Rev. Lett. 94 (2005) 171301, astro-ph/0409403.
- [24] E. Komatsu, et al., WMAP Collaboration, arXiv: 0803.0547.
- [25] J.F. Navarro, C.S. Frenk, S.D.M. White, Astrophys. J. 462 (1996) 563, astro-ph/9508025.
- [26] L. Bergstrom, P. Ullio, J.H. Buckley, Astropart. Phys. 9 (1998) 137, astro-ph/9712318.
- [27] C.-N. Yang, Phys. Rev. 77 (1950) 242.
- [28] K. Hagiwara, R.D. Peccei, D. Zeppenfeld, K. Hikasa, Nucl. Phys. B 282 (1987) 253.
- [29] D. Maurin, F. Donato, R. Taillet, P. Salati, Astrophys. J. 555 (2001) 585, astro-ph/0101231.
- [30] I.V. Moskalenko, A.W. Strong, Astrophys. J. 493 (1998) 694, astro-ph/9710124.
- [31] E.A. Baltz, J. Edsjo, Phys. Rev. D 59 (1999) 023511, astro-ph/9808243.
- [32] E.A. Baltz, J. Edsjo, K. Freese, P. Gondolo, Phys. Rev. D 65 (2002) 063511, astro-ph/0109318.
- [33] O. Adriani, et al., arXiv: 0810.4995, 2008.
- [34] K.R. Dienes, C.F. Kolda, J. March-Russell, Nucl. Phys. B 492 (1997) 104, hep-ph/9610479.
- [35] S.A. Abel, M.D. Goodsell, J. Jaeckel, V.V. Khoze, A. Ringwald, JHEP 0807 (2008) 124, arXiv: 0803.1449.
- [36] M. Endo, M. Kawasaki, F. Takahashi, T.T. Yanagida, Phys. Lett. B 642 (2006) 518, hep-ph/0607170.
- [37] M. Endo, F. Takahashi, T.T. Yanagida, Phys. Lett. B 658 (2008) 236, hep-ph/0701042.
- [38] M. Endo, F. Takahashi, T.T. Yanagida, Phys. Rev. D 76 (2007) 083509, 0706.0986.
- [39] Y.B. Zeldovich, I.Y. Kobzarev, L.B. Okun, Zh. Eksp. Teor. Fiz. 67 (1974) 3.
- [40] T.W.B. Kibble, J. Phys. A 9 (1976) 1387.
- [41] A. Vilenkin, Phys. Rev. D 23 (1981) 852.
- [42] D. Coulson, Z. Lalak, B.A. Ovrut, Phys. Rev. D 53 (1996) 4237.
- [43] S.E. Larsson, S. Sarkar, P.L. White, Phys. Rev. D 55 (1997) 5129, hep-ph/9608319.
- [44] F. Takahashi, T.T. Yanagida, K. Yonekura, Phys. Lett. B 664 (2008) 194, arXiv: 0802.4335.

⁴ The significant amount of the gravitational waves may be produced in the collisions of domain walls [44]. The hidden gauge bosons are also produced by the annihilation processes of the domain walls.

⁵ There also exist antiprotons produced from the A' decay, whose flux can also be calculated and compared with the present data. However, we did not pursue the detailed comparison in this Letter due to the large experimental uncertainties and our poor understanding of diffusion models.